

**UNTERSUCHUNG NICHTLINEARER VORGÄNGE
BEI DER STIMULIERTEN BRILLOUINSTREUUNG
BEI EINEM AUSGEDEHNTEN PLASMA
UNTERKRITISCHER DICHT**

DISSERTATION

zur Erlangung des Grades eines
Doktors der Naturwissenschaften
der Abteilung für Physik und Astronomie
an der Ruhr-Universität Bochum

von

Syed Abbas Hussain Rizvi

aus

Islamabad (Pakistan)

Bochum 1983





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INVESTIGATION OF
NONLINEAR PROCESSES ASSOCIATED WITH STIMULATED BRILLOUIN
SCATTERING IN AN UNDERDENSE AND EXTENDED PLASMA

by

Syed Abbas Hussain Rizvi

This work is submitted as a dissertation in partial
fulfillment of the requirements of the degree
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- Contents

	Page	
List of Diagrams	III	
Deutsche Zusammenfassung	VIII	
CHAPTER 1	Introduction	1
CHAPTER 1	Theoretical Background	5
2.1	Theory of parametric decay processes	5
2.1.1	Physical mechanism	6
2.1.2	Linear theory	8
2.1.3	Non-linear theory	17
2.2	Saturation mechanisms	23
2.2.1	Steepening of the density profile	23
2.2.2	Rescattering	23
2.2.3	Secondary decay of ion acoustic wave	24
2.2.4	Self focusing and filamentation	25
2.2.5	Ion trapping	27
2.2.5.1	Threshold of ion trapping	37
2.2.5.2	Side bands	39
2.2.5.3	Effect of collisions	40
2.2.6	Ion acoustic wave harmonics	42
2.2.7	Wave breaking	43
2.2.8	Ion heating	44
2.2.9	Pump depletion	46
CHAPTER 3	Experimental Set-up	47
3.1	The high power CO ₂ laser system	47
3.1.1	Optimization of the laser system	58
3.1.2	Characteristics of the high power laser system	61
3.2	The Z-pinch	64
3.3	The scattering experiment	66
3.3.1	The forward scattering experiment	67
3.3.2	The backward scattering experiment	71

	page
CHAPTER 4	Results and Discussion 76
4.1	Parametric decay processes 76
4.1.1	Stimulated Raman scattering and Stimulated Compton scattering 76
4.2	Stimulated Brillouin scattering 77
4.2.1	Back scattered radiation 78
4.2.2	Time integrated back scattered light spectrum 81
4.2.3	Forward scattering 85
4.3	Time resolved back scattered light spectrum 89
4.4	Split spectra 97
CHAPTER 5	Comparison 111
CHAPTER 6	Summary 120
REFERENCES	128
APPENDIX	134
A.1	Specifications 134
A.2	Basic CO ₂ laser theory 138
A.3	Numerical modelling of the CO ₂ laser 140
A.4	E-beam resonator 144
A.5	Physical mechanism of the e-beam laser 149
A.6	De-excitation processes 154
A.7	The computer model 1 156
A.8	Comparison of computer results with experiment 164
A.9	The computer model 2 180
A.10	List of symbols 185
REFERENCES	190

List of Diagrams

	Page
Fig.(2.1) A schematic diagram showing regions of different types of three wave interaction	5
Fig.(2.3) A schematic diagram of the density modulation, associated currents and magnetic field, and the directions of the forces responsible for stimulated scattering	7
Fig.(2.4) An (ω, k) diagram for SBS and SRS	10
Fig.(2.5) Wave vector diagram for scattering	14
Fig.(2.6) Physical picture for the production of filamentation	26
Fig.(2.7) Electron temperature as a function of laser intensity	30
Fig.(2.8) Density modulation as a function of T_e/T_i	31
Fig.(2.9) Bounce frequency as a function of density modulation	32
Fig.(2.10) Trapped particle orbits in ion-acoustic wave along with the corresponding wave potential and ion velocity distribution $f(v_i)$	34
Fig.(2.11) g is plotted as a function of the density modulation	39
Fig.(2.12) H is plotted as a function of the density modulation	41
Fig.(2.13) Sequence of events leading to wave breaking	42
Fig.(2.14) Ion temperature as a function of RI_0 as predicted by the ion heating model	45

	Page
Fig.(3.1) A schematic diagram of the high power CO ₂ laser system	48
Fig.(3.2) Modulated and the injection mode locked output of the helical TEA CO ₂ laser along with the near field distribution	50
Fig.(3.3) Helical TEA CO ₂ laser output with mode locking	51
Fig.(3.4) A schematic diagram of the e-beam laser system	52
Fig.(3.5) Annular field distribution of the e-beam laser	54
Fig.(3.6) Far field energy distribution of the e-beam laser	55
Fig.(3.7) E-beam laser output	56
Fig.(3.8) The e-beam laser output power as a function of the anode-cathode spacing	59
Fig.(3.9a) E-beam laser peak output power as a function of sustainer voltage for different gas ratios	60
Fig.(3.9b) E-beam laser output energy as a function of sustainer voltage for the optimum gas ratio	61
Fig.(3.10) A schematic diagram showing the process of injection mode locking	62
Fig.(3.11) A schematic diagram of the Z-pinch device	65
Fig.(3.12) Radial density distribution of the plasma column in Z-pinch	66

	Page
Fig.(3.13) A schematic experimental arrangement for the investigation of temporal and spectral aspects of SBS	68
Fig.(3.14) Bremsstrahlung emission signal of the plasma as a function of time	71
Fig.(3.15) Experimental arrangement for the temporal resolution of the back scattered light spectra	73
Fig.(3.16) Incident and the back scattered pulses	74
Fig.(4.1) Relative back scattering power as a function of laser intensity in several shots	79
Fig.(4.2) Relative back scattering power as a function of laser intensity within a shot	80
Fig.(4.3) Time integrated spectrum of the back scattered and the incident laser light	82
Fig.(4.4) Frequency shift in the back scattered radiation relative to the incident laser radiation as a function of the time coincidence	83
Fig.(4.5) Time integrated spectrum of the incident laser light (P 20) and the scattered light (FS) in the forward direction	87
Fig.(4.6) Oscilloscopic trace of the incident laser light and the FS light	88
Fig.(4.7) Back scattered laser pulses and the spectrally resolved signals	90
Fig.(4.8) Time resolved back scattered light spectra of three laser pulses in a shot	92

	Page
Fig.(4.9) A time resolved back scattered laser pulse spectra	93
Fig.(4.10) A typical non-splitted spectra of a back scattered laser pulse	94
Fig.(4.11) The time integrated and time resolved spectra	95
Fig.(4.12) The ratio between the frequency shift of the two back scattered pulses is plotted as a function of the ratio between the intensities of the corresponding laser pulses	96
Fig.(4.13) Spectral resolution of a back scattered laser pulse	98
Fig.(4.14) Time resolved five different wavelength channel signals corresponding to a back scattered laser pulse	99
Fig.(4.15) Frequency difference between the two peaks in a splitted spectrum as a function of the back scattering	101
Fig.(4.16) Temporal evolution of a back scattered laser pulse	103
Fig.(4.17) Time coincidence between an incident and a corresponding back scattered laser pulse	103
Fig.(4.18) SBS as a function of laser intensity for three laser pulses along with their time resolved spectra	106
Fig.(4.19) An envelope function with an oscillating function and its Fourier transform	109
Fig.(4.20) Density modulation as a function of SBS	109

	Page
Fig.(5.1) Experimental measurements of back scattering versus laser intensity in case of gas and double laser pulse solid target as reported in the literature	112
Fig.(5.2) E-beam laser output without IML and with IML by a cw-laser along with the corresponding back scattered signals	114
Fig.(5.3) Oscilloscope traces of incident and back scattered laser pulses near threshold and at large back scattering	115

Deutsche Zusammenfassung

Einführung

Im Falle der Laserfusion scheint das Erfordernis für eine wirksame Laserlichtabsorption im Target noch durch einige nichtlineare Prozesse, die das Laserlicht streuen, beeinträchtigt zu werden. Diese können einen bedeutenden Teil der Laserstrahlung zurückstreuen. Dabei ist die Stimulierte Brillouin Streuung (SBS) Gegenstand vieler Untersuchungen gewesen, weil bei diesem Prozeß entsprechend der Manley-Rowe-Beziehung fast 100% der vom Laser gelieferten Strahlungsenergie vom Plasma zurückgestreut werden kann.

Jedoch eignen sich lasererzeugte Pellet-Plasmen nicht für einen Vergleich mit der Theorie, weil verschiedene Nachteile damit verbunden sind. Z.B. ist es schwierig, die Bedingungen für die Wechselwirkung zu kennen, solange das Plasma vom selben Laserpuls sowohl erzeugt als auch untersucht wird, weil das Plasma sehr starke Dichte- und Temperaturgradienten hat. Sogar sogenannte experimentelle Modellaserplasmawechselwirkungsergebnisse von unterkritisch dichten und mehr oder weniger gut vorionisierten Plasmen zeigen eine weite Streuung, z.B. liegt der Sättigungsgrad für SBS zwischen 0,3% /84/ und 85% /78/.

Um die unterschiedlichen Ergebnisse zu verstehen, die in der Literatur genannt werden, ist ein Laserwechselwirkungsexperiment durchgeführt worden. Mit diesem Experiment sollte der SBS-Prozeß unter wohl definierten Bedingungen untersucht werden.

In der ersten Arbeit /16,31/ konnte durch Versuche an einem wohl definierten Plasma zu einem besseren Verständnis dieser weiten Streuung der bisherigen Meßergebnisse beigetragen werden. Viel konnte mit Hilfe des entwickelten Modells der Plasmabeeinflussung durch die Laserbestrahlung ge-

klärt werden. Aus diesem Modell läßt sich T_e/T_i bestimmen, weil die Rotverschiebung $\Delta\nu$ der SBS-Linie mit der Summe von T_e und T_i zusammenhängt. So ließen sich T_e und T_i im Wechselwirkungsvolumen im einzelnen bestimmen. Damit ist eine Bestimmung des theoretischen Schwellwerts und der Anwachsraten der SBS möglich. Diese theoretische Berechnung stimmt mit den experimentellen Ergebnissen überein und damit ist die Anwendbarkeit der theoretischen Beschreibung von Forslund /8/ auf die Bedingungen unseres Plasmas erwiesen.

In seiner Dissertation ist Handke /31/ ein recht befriedigendes Verständnis der SBS-Vorgänge in Schwellwertnähe gelungen. Die aus dem Heizmodell vorhergesagte Abhängigkeit der Rotverschiebung $(P/A)^{1/7,5} \propto \Delta\nu$ hat sich zwar auch bei einzelnen SBS-Spektren weit oberhalb des Schwellwerts erkennen lassen. Doch der aus der linearen SBS-Theorie erwartete steile Anstieg der Rückstreuung trifft nur bis zu einer relativen Rückstreuleistung von 10% zu. Ab 10% setzt ein zusätzlicher Dämpfungsmechanismus den weiteren Anstieg der Rückstreuung mit der Bestrahlungsstärke stark herab. Weil dieser Dämpfungsprozeß erst mit ansteigender Wellenamplitude einsetzt, muß es sich wohl um einen nichtlinearen Prozeß handeln.

Es war die zentrale Aufgabe der vorliegenden Arbeit, diesen Mechanismus zu identifizieren und zu klären, der sich auf das Verständnis in der SBS bei einer Bestrahlungsstärke oberhalb des Schwellwertbereichs konzentrierte. Für das vollständige Verständnis des rückgestreuten Lichtes mußten nicht nur die SBS, sondern auch die Stimulierte Raman Streuung (SRS) und die Stimulierte Compton Streuung (SCS) berücksichtigt werden.

Theoretischer Hintergrund

Für das Grundverständnis des Gegenstandes werden sowohl die lineare als auch die nichtlineare Theorie der Laserplasmawechselwirkung kurz erläutert. Es wird angenommen, daß das Plasma homogen und unterdicht ist. Zwischen vielen parametrischen Prozessen werden nur Stimulierte Brillouin-Streuung (SBS), Stimulierte Compton-Streuung (SCS) und Stimulierte Raman-Streuung (SRS) betrachtet, weil sie die Laserstrahlung in bedeutendem Umfang streuen können.

Aus der linearen Theorie wurden der Schwellwert und die Anwachsrate von SBS, SCS und SRS bestimmt, die dabei behilflich sind, den parametrischen Prozeß zu identifizieren, der für die Laserlichtstreuung verantwortlich ist. Da SBS am wichtigsten von ihnen ist, wurden die weiteren Untersuchungen des Gegenstandes auf SBS beschränkt. Es ist auch gezeigt worden, daß die Maximalstreuung in der Rückwärtsrichtung vorkommt, d.h. in der entgegengesetzten Richtung von $\vec{k}_0$, dem Ausbreitungsvektor der CO_2 -Laserstrahlung.

Auf der anderen Seite sagt die nichtlineare Theorie die wichtigsten Charakteristika des Streuungsprozesses voraus. Das wichtigste Charakteristikum ist die Sättigung der SBS. Es wurde herausgefunden, daß die Plasmawechselwirkungslänge, die Plasmadichte und die Ionendichtemodulationsamplitude den Umfang der zurückgestreuten Energie bestimmen.

Ein Versuch, unsere experimentellen Ergebnisse aufgrund der existierenden theoretischen Modelle zu erklären, ist fehlgeschlagen. Es wurde herausgefunden, daß diese Modelle nicht allgemeingültig sind.

Obwohl sehr starke SBS beobachtet worden ist, impliziert sie nicht eine vollständige Abwesenheit aller Sättigungsmechanismen. Diese Sättigungsmechanismen können das Anwachsen der SBS drastisch beeinflussen. Später wurde die Sättigungsmechanismenpräsenz experimentell bestätigt. Es wurde

festgestellt, daß die Stärke der Rückstreuung nicht so schnell anwächst wie es aufgrund der nichtlinearen Theorie erwartet werden konnte.

Eine theoretische Untersuchung der Sättigungsmechanismen hat gezeigt, daß nur Vielfachstreuung, Harmonischenerzeugung und Ioneneinfang in unserem Experiment wirksam sein können. Da eine Harmonischenerzeugung bereits von Gellert /13/ analysiert worden ist und die Existenz der Vielfachstreuung in der vorliegenden Arbeit nicht experimentell festgestellt werden konnte, wurde der Ioneneinfang im einzelnen untersucht. Die Dichtemodulation wurde als eine Funktion von T_e/T_i bestimmt, wobei T_e/T_i von der Laserintensität abhängt. T_e/T_i kann für eine bestimmte Laserintensität nach dem Heizmodell von Handke /31/ explizit berechnet werden. Die Kenntnis der Dichtemodulation war für den Ioneneinfang notwendig, um die "bounce"-Frequenz der Ionen zu schätzen. Die Schwellwertbedingung für den Ioneneinfang, die von der "bounce"-Frequenz abhängt, wurde festgestellt. Es wurde herausgefunden, daß in unserem Experiment eine Minimumdichtemodulation von 8% für den Ioneneinfang erforderlich ist. Es ist auch theoretisch erwiesen, daß der Ioneneinfang eine erhöhte Dämpfung und eine Erzeugung von "side bands" zur Folge hat. Die "side bands" werden aufgrund der Wechselwirkung zwischen der modulierten Ionen-akustischen Welle und den gefangenen Teilchen, die in dem Potentialtopf der Welle hin- und herlaufen, erzeugt. Zusätzlich ist auch die Wirkung der Stöße berücksichtigt worden, weil unser Plasma nicht stoßfrei war. Die Theorie für Elektroneneinfänge von Mima und Nishikawa /56/ wurde für den Fall der Ioneneinfänge durch Einschluß der Ionendynamik erweitert. In dieser Analyse sind alle Stoßtypen berücksichtigt worden. Eine grobe Schätzung hat gezeigt, daß in unserem Fall die Stöße für eine Dichtemodulation von 9% oder weniger wirksam werden können.

Experimenteller Aufbau

Zur Untersuchung parametrischer Prozesse bei der Laser-plasmawechselwirkung werden ein voll diagnostiziertes, fremd-erzeugtes Plasma und ein Laser mit einer Strahlungsstärke von mehr als 10^{15} Wm^{-2} , was einem Schwellwert von SBS entspricht, benötigt.

Das Hochleistungslasersystem

Ein Hochleistungs-CO₂-Lasersystem mit bestimmten spezifischen Erfordernissen ist entwickelt worden. Eine sehr kurz-brennweitige Fokussierungslinse wird nicht verwendet, um hohe Bestrahlungsstärken zu erzielen, weil dann die räumliche Ausdehnung des Laserfokus¹ sehr gering wird und damit auch das Wechselwirkungsvolumen kleiner wird.

Für ein sauberes Laserplasmawechselwirkungsexperiment ist ein Laseroutput mit einem hohen Kontrastverhältnis, einer einzelnen longitudinalen und transversalen Mode und einer kleinen Strahlungsdivergenz erforderlich. Außerdem sollte der Laserpuls nicht größer als 2 ns sein, um eine Dichtevariation im Wechselwirkungsvolumen des Plasmas während des Experiments zu vermeiden. Weiterhin sollte(n) der (die) Laserpuls(e) nicht moduliert werden, da sonst die Ergebnisse nicht reproduzierbar sind. Um eine Filamentierungsinstabilität, die mit SBS konkurrieren könnte, zu vermeiden, sollte der Laseroutput frei von heißen Flecken sein. Es sind viele Anstrengungen unternommen worden, um eine solche unstrukturierte Focalverteilung zu erreichen.

Alle oben genannten hohen Anforderungen an die Laserstrahlung werden zusammen mit dem Injection Mode Locking (IML) erfüllt, wie es im einzelnen im Anhang beschrieben wird. Es beruht auf einem Prinzip, bei dem ein kleiner Laseroszillator, der Master Oszillator genannt wird, dazu gebracht wird, eine Strahlung mit den gewünschten Eigenschaften zu erzeugen.

Diese Strahlung wird in einem großen Laser, der Slave Oszillator genannt wird, injiziert, wo dann die hohe Leistungsverstärkung mit den gewünschten Eigenschaften erfolgt. Als Master Oszillator wurde ein TEA-Wendellaser des semikonfokalen Typs angewandt. Dieser sendet einige hundert Kilowatt mit einer Halbwertszeit von 200 ns aus. Die Emission einer einzelnen longitudinalen und transversalen Mode wurde durch die Einführung einer Blende im Resonator und durch Austausch des Frontspiegels gegen einen Resonanzreflektor erreicht. Zusätzlich wurde ein Pulsschneidesystem angewandt, um einen 200 ns Puls auf einen Puls von einigen Nanosekunden Dauer zu verkürzen.

Für IML wurde der Output des Master Oszillators in den Slave Oszillator geschickt. Der Slave Oszillator war ein elektronenstrahlangeregtes (e-beam) CO_2 -Lasersystem mit einer instabilen Resonatorkonfiguration. Das Entladungsgefäß war mit einer gasförmigen Mischung von He und CO_2 mit einem Druck von 2×10^5 Pa gefüllt, und das aktive Volumen betrug ca. 10 l. Das N_2 -Gas war absichtlich vermieden worden, um die Energie der verzögerten Pulse, die die optischen Komponenten beschädigen können, zu reduzieren.

Das e-beam-Lasersystem wurde optimiert und die Charakteristika des Outputs wurden untersucht. Ein numerisches Modell wurde entwickelt, um die Laser-Output-Charakteristika und die Wirkung der verschiedenen Parameter darauf vorherzusagen. Es ist im Anhang gezeigt worden, daß die Computerergebnisse mit den experimentellen Ergebnissen gut übereinstimmen. Dadurch wird bestätigt, daß der Laser mit optimaler Leistung in Betrieb war und gleichzeitig sein Output alle Erfordernisse für ein Laserplasmawechselwirkungsexperiment erfüllte. Hiermit wird es ermöglicht, eine Leistungsdichte bis zu 10^{17} Wm^{-2} in der Brennebene der 714 mm Linse zu erreichen. Der Laser emittiert einen Zug von Nanosekundenpulsen mit einer Zwischenzeit von 16,7 ns, die genauso lang ist wie die "round trip"-Zeit des Resonators. In einem Laserplasmawechselwirkungs-

experiment ist ein Zug von Pulsen in dem Sinne von Vorteil, daß er mehrere Messungen in einem Schuß ermöglicht. Wenn der Zeitunterschied zwischen den Pulsen ausreicht, um die Gleichgewichtsbedingungen im Plasma wiederherzustellen, dann ist es gerechtfertigt, alle Laserpulse in einem Schuß als individuelle Laserpulse mit unterschiedlichen Peakleistungen zu betrachten. Es konnte gezeigt werden, daß diese Annahme gerechtfertigt war.

Z-Pinch

Auf der anderen Seite wurde das Plasma in einem dynamischen Z-Pinch mit einer Gesamtentladungsenergie von 20 KJ produziert. Dieses Plasma ist von Finken et al /69-71/ untersucht worden. Daraus ergeben sich wichtige Informationen über die Plasmamparameter, die für das Verständnis eines Laserplasmawechselwirkungsexperiments erforderlich sind. Z.B. gilt $K_B T_e \approx K_B T_i = 10 \text{ eV} \pm 10\%$, $Z_{\text{eff}} \approx 1$, und das Plasma bleibt für 80 ns um die Phase der Maximalkompression herum stationär, und es ist ziemlich homogen. Der größte Teil des Experiments wurde bei einer Plasmadichte von $0,16 n_{\text{cr}}$ durchgeführt, weil ein unterdichtetes Plasma erforderlich war und die Plasmamparameter für diese Bedingungen bekannt waren.

Die Plasmadichte in unserem Experiment wurde ganz unterdicht gehalten, obwohl die Möglichkeit besteht, nahe oder oberhalb der kritischen Dichte (n_{cr}) ein Plasma zu produzieren. Die Absicht war, ähnliche Bedingungen zu untersuchen, wie es sie in dem großen ausgedehnten Koronabereich der Laser-Pellet-Wechselwirkung gibt.

Wechselwirkungsexperimente

Zur Untersuchung der Laserplasmawechselwirkung ist der Output des Hochleistungs- CO_2 -Lasersystems in das Z-Pinchplasma fokussiert worden.

Verschiedene photo- und pyro-elektrische Detektoren sind angewandt worden, um den zeitlichen Verlauf der einfallenden Laserstrahlung, der vom Plasma rückgestreuten Strahlung und der Plasmaentwicklung zu registrieren. Gleichzeitig sind die einfallende und die rückgestreute Laserstrahlung durch ein Infrarotkamarasystem (IRSA) analysiert worden.

Um die dynamischen Auswirkungen zu vermeiden, wird der Laser mit dem Plasma so synchronisiert, daß er mit dem Plasma zur Zeit der Maximalkompression des Plasmas in Wechselwirkung steht.

Für die zeitliche Auflösung der Rückstreuenspektren war ein Array-Detektor hinter dem IRSA-System aufgebaut.

Experimentelle Ergebnisse

Stimulierte Raman Streuung (SRS) und Stimulierte Compton Streuung (SCS)

Es wurde ein Versuch unternommen, um SRS bei einer niedrigen Plasmadichte von $0,08 n_{cr}$ zu beobachten, die notwendig ist, um die SRS-Wellenlänge zu dem Wellenlängenbereich, der beobachtet werden soll, zu verschieben. Trotz aller Anstrengungen wurde kein SRS Signal festgestellt. Auf der anderen Seite konnte SCS nicht gemessen werden, obwohl die Laserintensität um einiges über dem SCS Schwellwert lag. Die Abwesenheit sowohl von SCS als auch von SRS konnte als Konsequenz des kurzen Laserpulses interpretiert werden. In diesem Fall überschreitet T_e/T_i sehr schnell den Wert eins. Dann ist die Landau-Dämpfung nicht mehr wirksam. Das ermöglicht das Anwachsen von SBS zu hohen Werten.

Stimulierte Brillouin Streuung (SBS)

Ein homogenes, vollständig ionisiertes Wasserstoffplasma von $0,16 n_{cr}$ wurde mit einem Laser von einer Intensität unterhalb von 10^{17} Wm^{-2} bestrahlt. Die SBS-Strahlung ist sowohl in

der Vorwärts- als auch in der Rückwärtsrichtung im einzelnen beobachtet worden. Das Experiment wurde mit einer hohen zeitlichen- und einer Frequenzauflösung von 200 ps und 0,6 GHz bzw. einer Nachweisgrenze von 0,003 des einfallenden Lichts durchgeführt. Das IRSA-System war eine Größenordnung mehr empfindlich. Trotzdem konnte eine Vielfachstrahlung weder in der Vorwärts- noch in der Rückwärtsrichtung festgestellt werden.

Rückgestreute Strahlung

Zuerst wurde die relative Rückstreuung für eine Laserintensität von weniger als 10^{17} Wm^{-2} gemessen. Die Rückstreuungsleistung wurde bis 100% beobachtet. Die Schwellwertlaserintensität stimmt mit der SBS-Theorie gut überein. Die rückgestreute Strahlung wurde sowohl zeitlich als auch spektral untersucht. Die zeitliche Analyse der rückgestreuten Laserpulse hat ein Zeitintervall von ca. 500 ps zwischen dem überschreitenden SBS-Schwellwert und der Erscheinung einer großen Rückstreuung gezeigt. Es ist schon theoretisch bekannt, daß die reziproke Anwachsrate von SBS nahe des Schwellwerts ca. 25 ps /16/ beträgt. Daraus folgt, daß SBS vom thermischen Niveau aus anwächst, weil SBS in diesem Zeitpunkt um etwa neun Größenordnungen ansteigt. Es ist gezeigt worden, daß die Selbstfokussierung und die Turbulenz auf SBS unter den untersuchten Bedingungen keinen Einfluß haben.

Zeitlich integrierte rückgestreute Lichtspektren

Auf der anderen Seite wurden viele zeitlich integrierte rückgestreute Lichtspektren untersucht. Daraus folgt, daß die rückgestreute Strahlung eine Rotverschiebung von ca. 17 GHz relativ zu der einfallenden Laserstrahlung hat. Es wurde festgestellt, daß die Verschiebung und die Breite des rückgestreuten Lichtspektrums mit der Rückstreuung anwachsen, wobei SBS

von der Laserintensität abhängt. Eine genauere Untersuchung des Spektrums hat ergeben, daß es die Streuung wegen eines parametrischen Prozesses und nicht als Reflexion am bewegten Plasma gibt. Es ist festgestellt worden, daß die rückgestreute Strahlung SBS ist. Die Frequenzverschiebung hat es ermöglicht, die Summe der Elektronen- und der Iontemperatur zu bestimmen. Auf der Basis dieser Ergebnisse hat Handke /31/ ein Heizmodell entwickelt. Das Heizmodell bestimmt explizit T_e und T_i und wurde später von Gellert experimentell bestätigt.

Vorwärtsstreuung

Außer der Rückwärtsstreuung wurde auch die Vorwärtsstreuung untersucht. Es wurde herausgefunden, daß es keine echte Vorwärtsstreuung gibt, sondern daß das frequenzverschobene Licht in Vorwärtsrichtung das rückgestreute SBS-Signal ist, das zu dem e-beam-Resonator zurückläuft, von ihm reflektiert wird und wieder das Plasma bestrahlt, wenn dieses transparent wird.

Trotz all dieser wertvollen Informationen ist es nicht möglich, die große Abweichung der rückgestreuten Kurve zu verstehen. Die Abweichung von den theoretischen Berechnungen ist in der Nähe von 10% SBS deutlich zu erkennen. Ein anderes Paradox war die Abhängigkeit der Frequenzverschiebung von der Laserintensität. Handke's Heizmodell /31/ sagt ein Anwachsen der Frequenzverschiebung mit der Laserintensität $I_0^{1/7,5} \propto \Delta\nu$ voraus, wohingegen die Theorie von Forslund /8/ $I_0^{1/3} \propto \Delta\nu$ im Falle einer starken Kopplung voraussagt. Auf der anderen Seite werden die Spektren in dem Maße breiter, in dem SBS anwächst. Das bewirkt, daß sich die Rückstreuungsspektren aller Pulse eines Schusses überlappen. Auf diese Weise ist es schwierig, die Gültigkeit des $I_0^{1/3}$ oder des $I_0^{1/7,5}$ Gesetzes in einem zeitlich integrierten rückgestreuten Lichtspektrum zu überprüfen.

Um die Art der Sättigungsmechanismen, die eine anomale Dämpfung nahe 10% SBS verursachen, und um die Abhängigkeit der Frequenzverschiebung von der Laserintensität zu überprüfen, wurde das rückgestreute Lichtspektrum zeitlich aufgelöst. Es wurde herausgefunden, daß die SBS-Spektren sich nahe 10% SBS in zwei Spitzen aufspalten und dabei eine anomale Sättigung einsetzt.

Gespaltene Spektren

Eine Untersuchung einer großen Zahl von gespaltenen und nichtgespaltenen Spektren hat gezeigt, daß die Frequenzverschiebung dem $I_0^{1/7,5}$ Gesetz folgt. Auf diese Weise ist Handke's Heizmodell /31/ bestätigt worden. Es wurde auch festgestellt, daß die gespaltenen Spektren aus zwei "side bands" um die Hauptwelle herum bestehen. Eine weitere Untersuchung des gespaltenen Spektrums mit zeitlicher Auflösung besser als 200 ps hat gezeigt, daß die Spektrumaufspaltung ($\delta\nu$) der zweifachen "bounce"-Frequenz entspricht. Außerdem wurde festgestellt, daß $\delta\nu$ allmählich mit dem Anwachsen der SBS abnimmt (im Bereich von 10-30% SBS). Eine solche Abnahme der Dichtemodulation bei anwachsender SBS ist auch von Gellert /13/ beobachtet worden.

Vergleich mit den Messungen anderer Autoren

Um unsere Experimente mit anderen zu vergleichen, wurde der Laser "injection mode locked", nicht "injection mode locked" und "injection locked" durch einen Dauerstrichlaser in Betrieb genommen. Dadurch konnte die Richtigkeit unserer Erklärung der Beobachtungen anderer Autoren bestätigt werden. Unser Modell erklärt so nicht nur unsere Ergebnisse, sondern auch viele unterschiedliche Ergebnisse, die in der Literatur berichtet worden sind. Es hat auch Hinweise darauf gegeben, warum es in unserem Experiment eine sehr hohe SBS gab, aber nicht in anderen Experimenten.

Außerdem wurde gezeigt, daß unsere Ergebnisse gut mit den kürzlichen Beobachtungen von Walsh und Baldis /54/ über die Ionenwellenamplitudenmodulation und die Harmonischenerzeugung übereinstimmen.

Zusammenfassung

Eine Untersuchung der Stimulierten Brillouin Streuung (SBS) ist unter wohl definierten Wechselwirkungsbedingungen durchgeführt worden. Die Rückstreuung bis fast 100% ist vom unterkritisch dichten Plasma mit schwachen Temperatur- und Dichtegradienten bei großer räumlicher Ausdehnung bei Bestrahlungsstärken des einfallenden Lasers von weniger als 10^{17} Wm^{-2} beobachtet worden. Die genannten Plasmabedingungen treten in ähnlicher Form in der Korona von Plasmen auf, weil sie bei der beabsichtigten Laserfusion vorkommen können.

Eine zeitliche Auflösung der rückgestreuten Spektren bestätigt das Heizmodell von Handke /31/ und gibt eine Erklärung für die große Abweichung der rückgestreuten Kurve von den theoretischen Berechnungen nahe bei 10% SBS.

Alle Möglichkeiten, die eine Spektrumsaufspaltung verursachen können, sind betrachtet worden. Daraus ergibt sich, daß die Spektrumsaufspaltung nur mit dem Ioneneinfang zusammenhängt. Zusammen mit dem Heizmodell von Handke /31/ und den Ergebnissen von Gellert /13/ ergibt sich folgendes Bild der Wechselwirkungsvorgänge:

Der scharfe Laserpuls läßt T_e/T_i sehr schnell über den Wert eins hinaus anwachsen, macht die Landau-Dämpfung unwirksam und läßt die Amplitude der Ionen-akustischen Welle anwachsen. Sobald die Ionenwellenamplitude genügend groß ist, daß 10% oder mehr Rückstreuung erfolgt, kann ein Teil der Ionen aus der Gesamtverteilungsfunktion eingefangen werden. Die eingefangenen Ionen werden auf Kosten der Ionenwellenenergie beschleunigt. Das verursacht eine starke Dämpfung für die Ionenwelle. Tatsächlich zeigt die Rückstreuungs-

kurve an diesem Punkt eine scharfe Krümmung, die den Anfang eines zusätzlichen Dämpfungsprozesses anzeigt. Auf diese Weise kann die Krümmung mit dem Anfang des Ioneneinfangs in Verbindung gebracht werden. Wegen der anwachsenden Dämpfung wird die Welle kleiner und setzt sich an einer niedrigeren Amplitude fest. Die Rückstreuung würde auch geringer werden, wenn die Wechselwirkungslänge nicht anwachsen würde (Gellert /13/). Aus Gellerts Arbeit ist bekannt, daß sich nach diesem Übergang die Dichtemodulation $\delta n/n$ bei Werten unterhalb von 10% in dieser quasistationären Phase mit anwachsender Wechselwirkungslänge festsetzt, um die gemessene Rückstreuung zu produzieren. In dieser folgenden quasistationären Phase wird entsprechend den Ergebnissen von Gellert /13/ die anomale Dämpfung der Harmonischenerzeugung bewirkt. Dies setzt voraus, daß vorher genügend Ionen über den beschriebenen Prozeß eingefangen worden sind und die sinusförmige Welle verzerren.

Nur während dieser Anfangsphase hat die Ionenwelle eine ausreichende Amplitude, um die "bouncing"-Oszillationen stark genug zu machen, so daß die Ionenwellenamplitude moduliert wird und die Aufspaltung des rückgestreuten Spektrums wie auch die erhöhte Dämpfung der Ionen-akustischen Welle bewirken kann.

Auch in anderer Hinsicht ergänzen die Ergebnisse dieser Arbeit die Resultate von Gellert und zeigen die Anwesenheit des Ioneneinfangs als einen wirksamen Sättigungsmechanismus. Er fand heraus, daß die Dichtemodulation auch mit dem Anwachsen der Laserintensität, die zu höherer Rückstreuung führt, kleiner wird, weil das entsprechende Anwachsen der Wechselwirkungslänge diese Abnahme überkompensiert. Das stimmt mit unseren experimentellen Ergebnissen überein, bei denen herausgefunden wurde, daß die "bounce"-Frequenz als ein Maß für die Dichtemodulation ansteigt. Die "bounce"-Frequenz ist direkt proportional zu der Wurzel aus der Dichtemodulation.

Wenn unser Modell und die experimentellen Ergebnisse von Gellert kombiniert werden, dann kann die Abhängigkeit der Dichtemodulation der SBS vorhergesagt werden. Es wurde festgestellt, daß zuerst die Dichtemodulation abrupt wegen der Ioneneinfangs abnimmt, wenn SBS über 10% ansteigt. Der Ioneneinfang führt zur Deformation der sinusförmigen Ionenwelle und auf diese Weise werden Harmonische produziert. Bei 30% SBS wird die Amplitude der Ionen genügend groß, um alle Ionen einzufangen. Deshalb wird die Ionenwelle nicht mehr gedämpft. Ein weiteres Anwachsen der Laserintensität steigert die Amplitude der Dichtemodulation, weil keine Ionen mehr übrig bleiben, die beschleunigt werden können.

Unser Modell erklärt nicht nur unsere Ergebnisse befriedigend, sondern trägt auch zum Verständnis anderer in der Literatur genannter Ergebnisse bei. So läßt sich z.B. die kürzlich von Walsh und Baldis /54/ beobachtete Modulation der SBS-Ionenwelle gut erklären und gestattet es, in Verbindung mit dem Heizmodell von Handke /31/ detaillierte Angaben über deren Wechselwirkungsverhältnisse zu machen.

Chapter 1

Introduction

Mankind is facing a severe energy crisis. World progress is threatened as the natural resources are inevitably being exhausted. The estimated world reserve of fossil fuels may be depleted within a century. One possible solution is Nuclear Fusion, which promises to be an inexhaustible source of energy.

Broadly, there are two major approaches to achieve controlled thermonuclear fusion reaction: magnetic confinement of hot plasmas at moderate densities, and inertial confinement at very high densities. Until 1970, all the research effort was focussed on magnetic confinement. In this approach long confinement time is required for the efficient burn of the low density reactants. Since 1970, another approach, which relies on inertial confinement, fusion in superdense plasma, has been in rapid development after the advent of high power pulsed lasers. It requires no magnetic field for the confinement, rather, it relies on inertia to compress the reactants.

One of the essential requirements for fusion is to heat the target to a high temperature. It imparts sufficient thermal velocity to the reactant nuclei for the latter to come close enough so that the short range nuclear forces can overcome the Coulomb barrier during collisions, such that fusion occurs. Thus, the reacting particles must be confined at high density and high temperature for a certain period of time. Laser beam can compress and heat a fuel pellet to very high densities (10^{32} ions m^{-3}) and very high temperatures (KeV), so that the reaction can proceed extremely rapidly in the form of a micro-explosion before the pellet disassembles (blows apart and cools).

Worldwide efforts have been made to produce very high power lasers. The novel properties of laser beams for localizing energy to $\sim 10^{-9}$ m^2 in area, field of $\sim 10^{16}$ $V\ m^{-1}$, and intensity up to $\sim 10^{25}$ $W\ m^{-2}$ are the bases for the effort.

The subject of Laser Plasma Interaction (LPI) has attracted a great deal of interest because of its possible application to laser plasma fusion. It is also being considered for the laser heating of magnetically confined plasmas /1/, as it provides the capability to simultaneously create thermonuclear temperatures and very high densities without the use of very high voltage low-inductance capacitor banks. This reduces the required confinement time.

Both approaches have a few inherent problems which can not yet be surmounted. In laser fusion approach, the requirement for efficient laser light absorption in target appears still to be threatened by some non-linear process which scatters the laser light.

The classical absorption of laser light by a plasma is due to electron-ion collisions. It is inefficient at low densities and high temperatures as the collision rate is directly proportional to the density-square and to $T^{-3/2}$. Enhanced absorption occurs only in the vicinity of the critical surface (resonance absorption). But the laser radiation has to penetrate through the subcritical region into the critical density region in order to satisfy the resonance condition between its frequency and the plasma frequency. Some parametric instabilities may occur in the subcritical density region. They may back scatter the laser radiation to a significant amount. In this way, the energy transport through the subcritical region is being threatened by the parametric instabilities and thus is a serious problem for laser fusion.

The back scattering can be due to various parametric processes. Among all these, Stimulated Brillouin Scattering (SBS) has been the subject of much theoretical /2-9,74/, numerical /8,11/, and experimental /75-86/ attention because of the possibility of near 100% back scattering.

To have the experimental confirmation of the theoretical prediction it would be quite important to know the

amount of laser energy backscattered due to SBS in a LPI experiment under different conditions. Comparing different measurements of the reflectivity of the laser irradiated targets, as reported in the literature, a very confusing picture emerges. In order to explain these widely different experimental results, various theoretical models have been put forward. Of course, they are not of universal nature. Even when the theoretical model is restricted to a particular experiment, the agreement between theory and experiment is usually rather poor.

In laser produced plasmas, it is difficult to establish the conditions for the interaction as long as the plasma is both produced and investigated by the same laser pulse, since the motion and inhomogeneities in the plasma complicate the problem to a great extent. In addition, there are formidable difficulties as it is not easy to control the output of a high power laser. In order to avoid such complications, the LPI experiment should be performed under well-defined conditions. Hence, the plasma should be preformed and well diagnosed, and the laser output must satisfy certain specific requirements, for instance, it should have single frequency and unmodulated output. To our knowledge, no such experiment has been reported. Therefore, an experiment was performed to study the physics of scattering process under well-defined conditions.

In our experiment, the plasma was not produced by a laser but was preformed by means of a Z-pinch. Although a plasma near or above the critical density could have been produced, our experiment was deliberately restricted to underdense plasma as (a) our interest was to investigate similar situations as encountered in the large extended coronal region of laser-pellet interaction which may endanger the prospects of laser fusion, (b) it is readily accessible to various diagnostic methods, and, (c) there is sufficient experimental data and theory available for comparison.

After a brief introduction of the subject, the theory of laser-induced parametric instabilities, as well as various saturation mechanisms are discussed in the second Chapter. The third Chapter contains the experimental details of the complete CO_2 laser system, the Z-pinch, and the laser plasma interaction experiment. Experimental results are reported in the next Chapter and are analyzed in the light of predictions discussed in the second Chapter. A comparison of our results with other experiments is made in Chapter 5, and the discrepancies are explained on the basis of a simple model. The last Chapter sums up the entire work. For a better understanding of the subject, the physics of the laser has been given in Appendix.

Chapter 2

Theoretical Background

The main emphasis of this thesis is to understand the nature of the scattered radiation due to SBS above threshold resulting from a laser plasma interaction experiment. First of all, some theoretical background is necessary for the understanding of the subject. This theoretical chapter is divided in four main sections. In the first section, the physical mechanism responsible for laser light scattering will be described, while in the second section a linear theory of LPI will be briefly reviewed. In the next section, a non-linear theory and some existing theoretical models will be briefly discussed and an attempt will be made to correlate them with our experimental results. Lastly, the effects of saturation mechanism on the backscattering will be described.

2.1 Theory of Parametric Decay Processes

When an intense laser light interacts with the plasma several Parametric Decay Processes (PDP) such as Stimulated Brillouin Scattering (SBS), Stimulated Raman Scattering (SRS), and Two Plasmon Decay (TPD) may occur. These processes occur in different density regions. In the experiment reported here, there was a provision to vary the target plasma density over a wide range (sections 3.3). A schematic diagram (Fig.2.1) shows the range of the plasma density in which a particular process may occur, if the intensity of the incident electromagnetic (em)-wave exceeds certain thresholds.

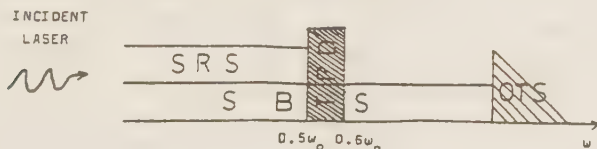


Fig.(2.1) Schematic diagram showing regions of different types of three wave interaction (in terms of plasma frequencies), when an em-wave of frequency ω_0 interacts with a homogeneous plasma.

SRS = Stimulated Raman Scattering

TPD = Two Plasmon Decay

SBS = Stimulated Brillouin Scattering

OTS = Oscillating Two-stream Instability

Among all the parametric decay processes, SBS and SRS may cause a serious problem to laser fusion. SBS involves a low frequency ion wave which takes up only a small amount of the incident laser light. Hence, it is theoretically possible (Manley-Rowe relation) to have almost 100% backscattering. Although SRS is less dangerous, it may nevertheless scatter a significant amount of the laser light and may create a hot electron environment which would make the achievement of a high density difficult. Hence, both processes will be discussed in the next section.

2.1.1 Physical Mechanism

Here, the physical mechanism responsible for the laser light scattering, for instance due to SBS, from an underdense plasma will be briefly described. Initially, the laser radiation (ω_o , k_o) will be incident on a plasma with thermal density fluctuation. This density fluctuation may scatter the incident em-wave. If this scattered em-wave (ω_- , k_-) leaves the plasma without affecting it, then the electrons will oscillate at frequencies ω_o and ω_- and nothing spectacular occurs. On the other hand, if the incident and the back scattered waves beat together, then not only the frequencies ω_o and ω_- but also $\omega_o + \omega_-$ and $\omega_o - \omega_-$ will be present in the plasma. The low frequency component will be more intense than the high frequency component since the pondermotive force varies as ω^{-2} . This pondermotive force is produced by the beating of the incident and the scattered em-waves, resulting in a low frequency force which can resonantly drive the ion mode at that frequency. In other words, the beating of the waves in the plasma will enhance the density fluctuation which scatters the incident laser light. This process can be described as a feedback loop (Fig. 2.2) which causes the ion wave and the backscattered wave to exponentiate together. The superscripts h(l) indicate the high (low) frequency component. These superscripts will be dropped later for the sake of simplicity.

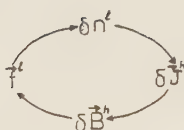


Fig.(2.2) A feedback loop showing the physical mechanism responsible for the exponentiation of the stimulated scattering instabilities. A density perturbation δn^l will produce an electron current $\delta \vec{J}^h$ resulting in a magnetic field of $\delta \vec{B}^h$. Electrons will experience a pondermotive force $\vec{r}^l$ and the process will exponentiate.

The sequence of events begins with a low frequency charge density perturbation δn^l . Directions of the forces $\vec{r}^l$ and currents $\delta \vec{J}^h$ responsible for the stimulated scattering of the incident em-wave are shown in Fig.(2.3).

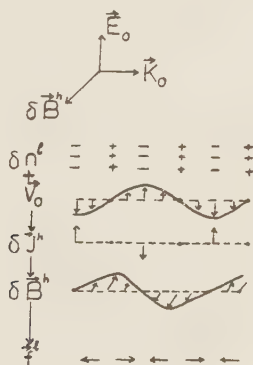


Fig.(2.3) Density modulation, associated currents and magnetic field, and the directions of the forces responsible for stimulated scattering are schematically shown here. Pump em-wave propagates horizontally. $\pm$ signs show the location of local extrema in the charge density modulations.

The charge density δn^1 together with the oscillatory velocity of the electrons $v_o^h = eE_o / (mw_o)$ in the high frequency pump wave with electric field E_o , produces a high frequency electron current

$$\delta \vec{j}^h = -ev_o^h \delta n_e^1 \quad (2.1)$$

This high frequency current produces a high frequency magnetic field $\delta \vec{B}^h$,

$$\mu_o^1 \nabla \times \delta \vec{B}^h = \delta \vec{j}^h + \epsilon_o \frac{\partial}{\partial t} \vec{E}^h \quad (2.2)$$

where $\epsilon_o (\mu_o)$ is the permittivity (permeability) of the vacuum. Due to this magnetic field, the plasma electrons feel a low frequency Lorentz force

$$\vec{f}^1 = -e(\vec{v}_o^h \times \delta \vec{B}^h) \quad (2.3)$$

Hence, the original density perturbation δn^1 grows still larger, and the process exponentiates.

2.1.2 Linear Theory

Here, an attempt will be made to present a linear theory of SBS and SRS within the frame work of the fluid model for a long homogeneous plasma. It is not a comprehensive review of the subject, rather only the aspects relevant to the experiment will be briefly discussed.

Dispersion relations for the natural modes in an unmagnetized plasma in the absence of a pump wave are:

$$w^2 = k^2 c^2 + w_{pe}^2 \quad (\text{em-wave}) \quad (2.4)$$

$$w^2 = 3k^2 v_e^2 + w_{pe}^2 \quad (\text{high frequency es-plasma wave}) \quad (2.5)$$

$$w^2 = k^2 c_s^2 \quad (\text{low frequency es-ion-acoustic wave}) \quad (2.6)$$

$$c_s^2 = \frac{\gamma_e Z K_B T_e}{m_i (1 + k^2 \lambda_{de}^2)} + \frac{\gamma_i Z K_B T_i}{m_i} \approx \frac{K_B T_e}{m_i} + \frac{3 K_B T_i}{m_i} \quad (2.6a)$$

$$v_j^2 = K_B T_j / m_j$$

$$\lambda_{de}^2 = \epsilon_o K_B T_e / (n_o e^2)$$

are ion-acoustic wave phase velocity, thermal velocity of jth type particles, and Debye length, respectively. In the above equations, w and k represent the frequency and wave number of a wave in the plasma, Z is the charge number, γ_j is the ratio of the specific heats c_p/c_v for j type

particles (j stands for electrons or ions), $\omega_{pe} = (n_0 e^2 / m_e \epsilon_0)^{1/2}$ is the electron plasma frequency, n_0 is the plasma density, T_j (m_j) is the temperature (mass) of the j type particle, and K_B is a Boltzmann-constant. Electrons move fast relative to the ion-acoustic wave, and hence can equalize their temperature rapidly. Thus, it has been assumed that $\gamma_e = 1$, and $\gamma_i = 3$.

Conservation of energy and momentum imposes the following matching conditions for all such processes (described above)

$$\begin{aligned}\omega_{\pm} &= \omega_0 \pm \omega \\ k_{\pm} &= k_0 \pm k\end{aligned}\quad (2.7)$$

Here, ω_0 is the frequency of the incident wave with wave number k_0 , $\omega_{\pm}$ ($k_{\pm}$) is the frequency (wavenumber) of the scattered wave, while $\omega(k)$ is the frequency (wavenumber) of the Langmuir or ion-acoustic wave. It can be best understood in an ω - k diagram, Fig.(2.4). It shows the dispersion curves of an em-wave, an ion-acoustic wave and an electron-plasma wave. The dispersion curves for the ion wave and electron plasma wave are basically horizontal lines compared to the parabola of the light wave dispersion (for small k) as ion wave velocity (c_s) is small compared to the plasma wave velocity (v_e) and both are much smaller than the velocity of light. Vectors ending on a particular dispersion curve represent that wave. The parallelogram shows that the matching conditions (2.7) are satisfied only for certain combination of waves, and then only a parametric decay process is possible. If the waves are not collinear, the parallelogram construction must be simultaneously possible in each orthogonal coordinate axis, thus making the matching conditions more difficult to satisfy.

Much effort has been invested in the theoretical work to understand the nature of the parametric decay processes. The first to deal with the possibility of a parametric instability was Silin / 2 /. His study was based on the hydrodynamic equations for a cold plasma neglecting pressure gradients, damping and coupling to the ion-acoustic wave. These factors were later included in Dubois et al Green's function analysis / 3 /. Lee et al / 4 / showed that the same results could

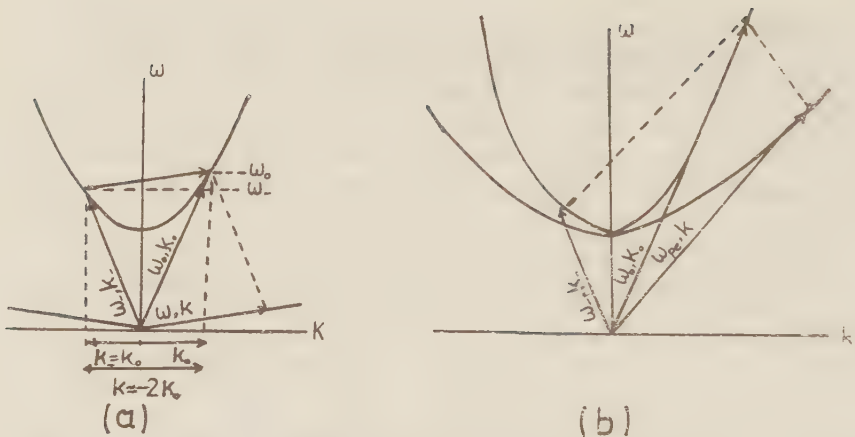


Fig.(2.4) An (ω, k) diagram for (a) SBS and (b) SRS. The straight lines are the dispersion relation for ion waves; the narrow parabolas, that for light waves and the wide parabolas, that for electron plasma waves. The vector like construction ensures that the matching conditions will be satisfied.

also be obtained from a fluid model with phenomenological damping constants. The kinetic theory based on the Vlasov equation was presented by Jackson / 5 /. He obtained the same results as before and also predicted the excitations of two unstable plasma modes with frequencies near the plasma frequency when allowance was made for the spatial variation of the pump field. Siebe Jörna / 6 / presented a linearized theory for the laser induced instabilities in a homogeneous plasma. He calculated the thresholds and growth rates for various instabilities. He extended the analysis to two dimensions and proved that SRS is primarily in the backward direction. Moreover, his analysis predicts a cutoff angle approaching 90° near a critical surface at one-fourth the critical density for SRS.

A systematic and comprehensive kinetic treatment of the interaction of an intense em-wave with a homogeneous plasma was presented by Drake et al / 7. /. In this analysis various instabilities were discussed. A linear and non-linear analysis of the one dimensional SRS and SRS in an infinite homogeneous plasma and a finite inhomogeneous plasma based on the fluid model was presented by Forslund et al / 8 /. They found that the non-linear fluid equations can predict backscatter energies of the order of the incident laser energy.

As our interest is to know the amount of incident laser light backscattered and in the conditions governing this back-scattering, the theory of Forslund et al, which is based on the fluid model, will be briefly reviewed. In this model, the following set of equations is used:

We define the electric potential φ and the vector potential $\vec{A}$ by $\vec{E} = -\frac{\partial \vec{A}}{\partial t} - \nabla \varphi$, $\vec{B} = \nabla \times \vec{A}$ (2.8)

where, $\partial/\partial t$ is the derivative. The wave equation is

$$c^2 \nabla^2 \vec{A} - \frac{\partial^2 \vec{A}}{\partial t^2} = -\epsilon_0^{-1} \vec{j} + \frac{\partial}{\partial t} (\nabla \varphi) \quad (2.9)$$

where, $\vec{j}$ is the current density.

The momentum equation is:

$$\frac{\partial \vec{v}_j}{\partial t} + (\vec{v}_j \cdot \nabla) \vec{v}_j = q_j m_j^{-1} (\vec{E} + \vec{v}_j \times \vec{B}) - \chi_j K_B T_j (m_j N_j)^{-1} \nabla N_j \quad (2.10)$$

where, q_j is the charge of species j , and N_j is its density.

The above equations (2.8) and (2.10) can be written as:

$$\begin{aligned} \frac{\partial \vec{v}_j}{\partial t} + (\vec{v}_j \cdot \nabla) \vec{v}_j &= q_j m_j^{-1} (-\nabla \varphi - \frac{\partial \vec{A}}{\partial t}) + q_j m_j^{-1} \vec{v}_j \times (\nabla \times \vec{A}) \\ &- \chi_j K_B T_j (m_j N_j)^{-1} \nabla N_j \end{aligned} \quad (2.11)$$

The continuity equation is

$$\frac{\partial}{\partial t} N_j + \nabla \cdot (N_j \vec{v}_j) = 0 \quad (2.12)$$

and the Poisson equation is

$$\nabla^2 \varphi = -\epsilon^{-1} \sum_j q_j N_j \quad (2.13)$$

To determine the quiver velocity of the particles in the em waves, we neglect thermal effects ($\gamma_j^{KT} \nabla N_j / m_j N_j = 0$) as well as the nonlinear terms $\{ (\vec{v}_j \cdot \nabla) \vec{v}_j, q_j m_j^{-1} \vec{v}_j \cdot \nabla (\nabla \times \vec{A}) \}$. Since the wave is purely em ($\varphi = 0$), (2.11) reduces to

$$\vec{v}_{j0} = -q_j m_j^{-1} \vec{A} \quad (2.14)$$

Here, it has been assumed that the waves propagate in the x-direction. The electron and ion densities can be separated into an homogeneous term and a spatially dependent perturbation as:

$$N_j = n_{0j} + \delta n_j(x) \quad (2.15)$$

Using (2.14) and (2.15), the wave equation (2.9) for electrons can be written as:

$$\left(\frac{\partial^2}{\partial t^2} + 2\Gamma_e \frac{\partial}{\partial t} + \omega_{pe}^2 - c^2 \nabla^2 \right) \vec{A} = -\omega_{pe}^2 \left(\frac{\delta n_e(x)}{n_0} \right) \vec{A} \quad (2.16)$$

where, $\delta n_e(x)$ is the fluctuating part of the electron density, Γ_e is the em-wave damping decrement, given by

$$\Gamma_e = \frac{\nu_{ei}}{2} \left(\omega_{pe} / \omega_0 \right)^2 \quad (2.17)$$

ν_{ei} is the electron-ion collision frequency.

To obtain a set of coupled equations for the incident and scattered wave vector potentials, the equations for the electron and ion density perturbation are required. If density modulations are not too large, then $\vec{v}_j$ in the nonlinear term may be neglected. Continuity equation (2.12) gives:

$$\nabla \cdot \vec{v}_j = -n_{0j}^{-1} \frac{\partial \delta N_j}{\partial t} \quad (2.18)$$

Taking the divergence of (2.10) and employing (2.18), (2.10) can be written as:

$$\begin{aligned} \frac{\partial^2 \delta n_j(x)}{\partial t^2} &= q_j m_j^{-1} n_{0j} \nabla^2 \varphi + (q_j / m_j)^2 (n_{0j} / 2) \nabla^2 \langle \vec{A} \cdot \vec{A} \rangle \\ &\quad + \nu_j v_j^2 \delta n_j(x) \end{aligned} \quad (2.19)$$

where, the following identities have been used:

$$\begin{aligned} \nabla (\nabla \times \vec{A}) &= \nabla (\nabla \cdot \vec{A}) - \nabla^2 \vec{A} = -\nabla^2 \vec{A} \\ \frac{\partial^2 N_j}{\partial t^2} &= \frac{\partial^2 \delta n_j(x)}{\partial t^2} \end{aligned}$$

For the electrons and ions, (2.19) can be explicitly written as:

$$\frac{\partial^2 \delta n_e(x)}{\partial t^2} = -(e/m_e) n_{oe} \nabla^2 \varphi + (e/m_e)^2 (n_{oe}/2) \nabla^2 < \vec{A} \cdot \vec{A} > + \gamma_e v_e^2 \nabla^2 \delta n_e(x)$$

$$\frac{\partial^2 \delta n_i(x)}{\partial t^2} = (e/m_i) n_{oi} \nabla^2 \varphi + \gamma_i v_i^2 \nabla^2 \delta n_i(x) \quad (2.20)$$

In the ion equation, the pondermotive effect of the incident em-wave has been neglected as it is smaller by a mass ratio. Elimination of $\nabla^2 \varphi$ from (2.20) yields:

$$\left(\frac{\partial^2}{\partial t^2} - c_s^2 \nabla^2 \right) \delta n(x) = n_{oj} (e^2/2m_e m_i) \nabla^2 < \vec{A} \cdot \vec{A} > \quad (2.21)$$

where, c_s is the ion-acoustic wave phase velocity as defined in (2.6) and $n_o \approx n_{oe} \approx n_{oi}$. As the ion-acoustic wave moves slowly, the quasi-neutral approximation is valid. A phenomenological damping term $2 \Gamma_i$ for the ion-acoustic wave in (2.21) may be added:

$$\left(\frac{\partial^2}{\partial t^2} + 2 \Gamma_i \frac{\partial}{\partial t} - c_s^2 \nabla^2 \right) \delta n(x) = n_{oj} (e^2/2m_e m_i) \nabla^2 < \vec{A} \cdot \vec{A} > \quad (2.22)$$

Consider the case of a linearly polarized laser wave, the vector potential for the pump can be written as:

$$\vec{A}_0(x, t) = -(E_0/\omega_0) e^{i(\omega_0 t - k_0 x)} \hat{e}_0 + c.c.$$

and,

$$E_0(x, t) = -\frac{\partial A_0}{\partial t} \quad (2.23)$$

Similarly, the backscattered wave and the ion-acoustic wave are given by:

$$\vec{A}_-(x, t) = A_- e^{i(\omega_- t - k_- x)} \hat{e}_- + c.c.$$

$$\delta n(x, t) = \delta n e^{i(\omega t - kx)} + c.c. \quad (2.24)$$

where, it has been assumed that the backscattered wave is also polarized in the same direction as the incident wave, and $\hat{e}_{0,-}$ is the unit vector in the direction of $\vec{A}_{0,-}$. As explained in section (2.21), E_0 will produce the density perturbation resulting in currents at $\omega \pm m\omega_0$, $k \pm k_0$, where m is an integer and $m=1$ corresponds to the lowest order coupling. Here, only the case of lowest order coupling will be considered, which will be later justified experimentally in Chapter (4).

The set of coupled equations (2.16) and (2.22) describes the parametric interaction of the em-wave and an ion-acoustic es-wave. If $A_0(A_-)$ is the vector potential of the incident

(scattered) em-radiation, then these differential equations describe the coupling between the incident and scattered waves and the plasma. They can be written by using (2.23) and (2.24) as:

$$\left(\frac{\partial^2}{\partial t^2} + 2\Gamma_i \frac{\partial}{\partial t} - c_s^2 \nabla^2 \right) \delta n(x) = n_{oe} (e^2 / 2m_e m_i) \nabla^2 \langle A_+ A_- \rangle \quad (2.25)$$

$$\left(\frac{\partial^2}{\partial t^2} + 2\Gamma_e \frac{\partial}{\partial t} + \omega_{pe}^2 - c^2 \nabla^2 \right) A_- = \omega_{pe}^2 (\delta n(x) / n_{oe}) A_+ \quad (2.26)$$

Pump depletion is neglected in this model. These two equations can be combined to yield the dispersion relation for SRS as:

$$(\omega^2 + 2i\omega\Gamma_i - k^2 c_s^2) =$$

$$0.5k^2 \omega_{pi}^2 v_o^2 \left[\omega^2 - 2\omega\omega_o - c^2(k^2 - 2kk_o \cos \vartheta) + 2i(\omega - \omega_o)\Gamma_e \right]^{-1} \quad (2.27)$$

where, ϑ is the scattering angle between $\vec{k}$ and $\vec{k}_o$, as illustrated in Fig.(2.5), and ω_{pi} is the ion-plasma frequency.

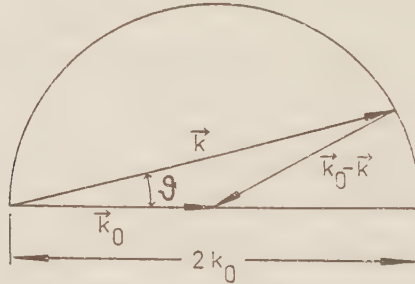


Fig.(2.5) Wave vector diagram for scattering at angles other than $\vartheta = \pi/2$ / 90° /. Here, $\vec{k}_o$, $\vec{k}_-$, $\vec{k}$ are the wave vectors of the incident, scattered and the ion-acoustic wave. Backscattering will occur when $\vartheta = 0$.

For the ion acoustic wave to grow, it is necessary that the denominator should vanish. Hence, the resonance condition is:

$$c^2(k^2 - 2kk_o \cos \vartheta) = \omega^2 - 2\omega\omega_o \quad (2.28)$$

which yields:

$$k = 2k_o \cos \vartheta \quad (2.29)$$

for $\omega = kc_s$. Here, ω^2 has been neglected in comparison with $2\omega\omega_o$, and c_s/c has been assumed much less than unity. k will be max-

imum when $\vartheta = 0$, which is the case for backscattering. In this case, a scattered Stokes component $\omega_-, \vec{k}_-$ propagates in the opposite direction to $\vec{k}_0$, as shown in Fig.(2.5).

In the case of forward scattering, ϑ will be $\pi/2$ and $k=0$ which leads to the conclusion that SRS in the forward direction is quite unlikely.

Stimulated Brillouin Scattering (SBS):

From the dispersion relation (2.27), threshold and the growth rate for SRS backscattering can be determined as:

$$(v_o/c)^2 = (\Gamma_e \Gamma_i \omega/\omega_c \omega_{pi}^2) \quad (2.30)$$

$$\gamma_{SBS} = (\omega_o \omega_{pi}^2 v_o^2 / \Gamma_i \omega c^2) \quad (2.30a)$$

where, the condition for the maximum value of k (2.29) has been used.

Stimulated Raman Scattering (SRS):

Similar to (2.25) and (2.26), the differential equations describing the coupling between the incident and scattered waves and the plasma for the case of SRS are:

$$\begin{aligned} & \left(\frac{\partial^2}{\partial t^2} + 2\Gamma_p \frac{\partial}{\partial t} - 3v_e^2 \nabla^2 + \omega_{pe}^2 \right) \delta n(x) \\ & = n_{oe} (e^2/2m_e^2) \nabla^2 (\vec{A}_0 \cdot \vec{A}_-) \end{aligned} \quad (2.31)$$

$$\begin{aligned} & \left(\frac{\partial^2}{\partial t^2} + 2\Gamma_e \frac{\partial}{\partial t} - c^2 \nabla^2 + \omega_{pe}^2 \right) A_- \\ & = -\omega_{pe}^2 (\delta n(x)/n_{oe}) A_0 \end{aligned} \quad (2.32)$$

where, Γ_p is the plasma wave damping decrement.

$$\Gamma_p = (\sqrt{\pi} \omega_{pe} / 2k^3 \lambda_{de}^3) \exp(-\frac{3}{4} - \frac{1}{2k^2 \lambda_{de}^2}) + \nu_{ei} \quad (2.33)$$

A Fourier transform of (2.31) and (2.32) yields the following dispersion relation for SRS as:

$$\begin{aligned} & (\omega^2 + 2i\omega\Gamma_p - \omega_{pe}^2 - 3k^2 v_e^2) = \\ & \omega_{pe}^2 k^2 v_o^2 [\omega^2 - 2\omega\omega_o - c^2(k^2 - 2kk_o \cos\vartheta) + 2i(\omega - \omega_o) \Gamma_e]^{-1} \end{aligned} \quad (2.34)$$

Again the resonance condition is:

$$c^2 (k^2 - 2kk_o \cos\vartheta) = \omega^2 - 2\omega\omega_o$$

$$\text{or } k = k_0 \cos \vartheta \pm (k_0^2 \cos^2 \vartheta - \frac{2\omega\omega_0}{c} + \frac{\omega^2}{c^2})^{1/2} \quad (2.35)$$

If ω^2 is neglected in comparison with $2\omega\omega_0$ in (2.35) then the maximum value of k occurs for $\cos \vartheta = (2\omega/\omega_0)^{1/2}$, and is:

$$k_{\max} = (2\omega_{pe}\omega_0)^{1/2} / c \approx 2k_0 \quad (2.36)$$

where, $\omega = \omega_{pe}$ for SRS. From the dispersion relation (2.34) and (2.36), the threshold and the growth rate for SRS will be:

$$(v_0/c)^2 = (\Gamma_p \Gamma_e / \omega_0 \omega_{pe}) \quad (2.37)$$

$$\gamma_{\text{SRS}} = (k^2 v_0^2 \omega_{pe} / 4\omega_0 \Gamma_p) \quad (2.37a)$$

The above mentioned results (2.30), (2.30a), (2.37), (2.37a) for the threshold and growth rate (just above threshold) of SRS and SRS are in good agreement with the kinetic treatment of C.S.Liu et al /9/.

Stimulated Compton Scattering (SCS):

When an intense laser beam is focussed on a plasma, then besides ordinary Compton collisions, in which a photon is absorbed and re-emitted spontaneously with a frequency shift in some given direction, the stimulated Compton effect must be taken into account. Stimulated Compton scattering will occur when an intense em-wave propagating in a plasma can be scattered not only off coherent plasma waves, such as electron plasma wave or ion-acoustic wave, but off electrons/ions bunches. Such a situation will arise, when $T_e \approx T_i$ resulting in the heavy damping of the ion-acoustic wave. The momentum exchange between the laser wave and the ions will take place via electrons. Further incoming photons will enhance this scattering process and a significant amount of scattering can be observed.

Stimulated Compton scattering both due to electrons and ions are possible. In case of ion Compton scattering, ions will scatter the incident laser wave and the phase velocity of the highly damped ion-acoustic wave will be almost equal to the ion thermal velocity. To treat this case, one must resort to a kinetic treatment of the low frequency quasi-mode.

As in the previous treatment, consider the interaction of an intense linear polarized em-wave with an homogeneous plasma, while allowing spatial variations in only one direction but allowing three components for velocity and fields. In such a situation, use of a linearized collisionless kinetic equation while including the pondermotive force will give the growth rate for SCS due to ions as/ 9 /:

$$\gamma_{SCS} = 0.5(v_o/v_e)^2 (\omega_{pe}^2/\omega_o) \text{Im}(1 + k_{de}^2 \chi_i)^{-1} \quad (2.38)$$

and the corresponding threshold is:

$$(v_o/c) = (v_e/c)(2\omega_o r_e/\omega_{pe}^2) \text{Im}(1 + k_{de}^2 \chi_i)^{-1/2} \quad (2.38a)$$

where, $\text{Im}\chi_i$ is the imaginary part of the ion susceptibility. Here, in the derivation of the equation (2.38), the ion-motion induced by the pondermotive force ($\vec{v} \times \vec{B}$) has been again neglected due to its heavy mass. An inclusion of this term will result in an additional term which is smaller than the original term by a factor of $(m_e/m_i)^2$.

2.1.3 Non-linear Theory

The linear theory of the laser plasma interaction as discussed in the previous section gives information about the threshold and growth rate of the parametric instabilities which helps to identify the scattering process. However, the most significant feature of a scattering process is its ability to scatter a certain amount of the incident energy. Therefore, in this section the three non-linear equations will be solved in order to find out the time development and spatial variation in the amplitude of the waves in the interaction volume.

As SCS grows above the threshold, non-linearities become important. They are, for example, the reaction of the back-scattered em-wave and the excited es-wave on the incident em-wave, described by the term $(\delta n/n_o)A_-$ in (2.16). There are many other non-linearities, like secondary decay of ion-wave and ion-trapping. They will be discussed in the next section. If higher order mode couplings (section 2.1.2)

and the damping of the light waves is neglected, then the amount of backscattering during a LPI experiment can be determined from the equations (2.16) and (2.22) of the previous section. They are:

$$\left(\frac{\partial^2}{\partial t^2} + \omega_{pe}^2 - c^2 \frac{\partial^2}{\partial x^2} \right) \vec{A}_0 = - \omega_{pe}^2 \left(\delta n_e / n_0 \right) \vec{A}_- \\ \left(\frac{\partial^2}{\partial t^2} + \omega_{pe}^2 - c^2 \frac{\partial^2}{\partial x^2} \right) \vec{A}_- = - \omega_{pe}^2 \left(\delta n_e / n_0 \right) \vec{A}_0 \quad (2.39)$$

$$\left(\frac{\partial^2}{\partial t^2} + 2\Gamma_i \frac{\partial}{\partial t} - c_s^2 \frac{\partial^2}{\partial x^2} \right) \delta n(x) = n_c (e^2 / 2m_e m_i) \frac{\partial^2}{\partial x^2} \langle \vec{A}_0 \cdot \vec{A}_- \rangle$$

where, the motion only in the x-direction has been considered.

Now, the case of a long homogeneous plasma will be considered. It will be assumed that the ion-density modulations associated with S3S have been driven to a certain value δn . The plasma has a constant density n_0 and is bounded within a slab of length L. It is known from (2.16) that the reaction of the (back) scattered em-wave and es-wave on the incident wave may be described by the term $(\delta n/n_0)\vec{A}$. Hence, the wave equations for both the incident and the scattered waves in stationary case under the assumption of the slowly varying amplitude can be written as:

$$\frac{\partial \vec{A}_-}{\partial x} = -\alpha' (\delta n/n_0) \vec{A}_0 \quad (2.40)$$

$$\frac{\partial \vec{A}_0}{\partial x} = -\alpha' (\delta n/n_0) \vec{A}_- \quad (2.41)$$

where,

$$\alpha' = (\pi/2) (n_0/n_{cr}) (1/\lambda_0) \left(1 - \frac{n_0}{n_{cr}}\right)^{-1/2} \quad (2.42)$$

λ_0 is the laser wavelength and n_{cr} is the critical density which is 10^{24} m^{-3} for CO_2 laser.

Normalizing the incident wave to unity ($A_0(0)=1$) and assuming that the initial backscatter noise level is small ($A_-(L)=0$), then the solution of equations (2.40) and (2.41) is:

$$A_0 = (e^{\beta'x} + e^{2\beta'L} e^{-\beta'x}) / (1 + e^{2\beta'L}) \quad (2.43)$$

$$A_- = (-e^{\beta'x} + e^{2\beta'L} e^{-\beta'x}) / (1 + e^{2\beta'L}) \quad (2.44)$$

where,

$$\beta' = \alpha' (\delta n / n_0)$$

Recognizing $|A_-(0)/A_0(0)|^2$ as the backscattered reflectivity, the above mentioned equations (2.43) and (2.44) yield / 10 /:

$$R = \left| \frac{e^{\beta'L} - e^{-\beta'L}}{e^{\beta'L} + e^{-\beta'L}} \right|^2 \quad (2.46)$$

$$= \tanh^2 \beta'L$$

It follows that the plasma interaction length, the plasma density, and the ion density modulation level determine the amount of the backscattering.

It must be pointed out here, that the theory thus far developed does not take into account the effect of the saturation mechanisms. It was assumed that the pump wave is quite large and unchanging, while the ion-acoustic wave and the backscattered waves are small and growing. Only the resonant interaction between the pump wave and one of the small amplitude waves were considered. However, when backscatter grows large enough some additional non-linearities which are generally known as saturation mechanisms may become important. These non-linearities will be discussed in the next section.

The amount of backscattering in a certain experiment could be explained either by the existing theoretical models or by equation (2.46) while incorporating the effect of the related saturation mechanisms. First of all some models will be briefly reviewed.

Many attempts have been made to understand the features of SBS. But early analysis / 2 - 6 / of SBS gives only the information about its threshold and growth rate. No attempt was made to predict the maximum amount of backscattering, ion density modulation and T_e/T_i . Generally, it was assumed that the instability exponentiates in space from some noise level, usually taken to be in the range of 10^{-2} — 10^{-4} .

Randall et al / 11 / explained the enhanced level of

scattering by considering the effects of a partially reflecting boundary. It was assumed that the laser is stationary but not the plasma. Hence, due to Doppler effect, the backscattered light can have the same frequency and will such serve as a source for SRS amplification. They discussed theoretically the growth of the backscattered SRS exclusively from the light reflected at the critical density. In this model, they assumed heavy damping for the ion-waves as a possible saturation mechanism. The model predicts that the reflecting surface can greatly increase the effective noise level, so that only a few growth lengths are required for significant scattering. Enhanced level of SRS was also expected even for stationary plasma.

Now, the validity of the above mentioned model will be checked for our experiment. In this model, it was assumed that the laser light reaches the critical surface and ion-heating is the possible saturation mechanism. Consequently, the model becomes invalid, if backscattering is strong as it will not allow a significant amount of the laser light to reach the critical surface. Moreover, it will be shown in section(2.2) that in our case ion heating is not a possible saturation mechanism.

Tajima / 12 / presented another model which describes the SRS and its saturation. He assumed that the produced ion waves get modulated and SRS can be saturated. Without specifying the process, he assumed that the ion-wave amplitude increases with the laser intensity resulting in an increased nonlinear modulations of ion waves. He derived an analytic expression for the reflectivity due to SRS with the help of a parametrically driven Burger's equation, Poisson and continuity equation for $k \lambda_{de} \ll 1$ as:

$$R = \Lambda (\delta n/n_0)^2 \quad (2.47)$$

where,

$$\eta = v_0/v_e$$

$$\Lambda = (w_0/w) (K_B T_e / m_e) \left\{ \frac{(k^2 + k_{de}^2)/k_{de}^2}{v_0 (w_0/w_{pe})} \right\}$$

$$k_{de} = \lambda_{de}^{-1}$$

In our case,

$$\begin{aligned}\eta &= 0.1 \text{ near threshold} \\ \eta &= 1.4 \text{ near } 100\% \text{ SBS} \\ \omega_{pe} &= 0.7 \times 10^{14} \text{ s}^{-1} \\ \omega_o &= 1.8 \times 10^{14} \text{ s}^{-1} \\ \omega &= 10^{11} \text{ s}^{-1}\end{aligned}$$

The validity of the Tajima model can be checked by the estimation of ion density modulation. For the above mentioned parameters, this model predicts an ion density modulation of $\sim 10^{-3}$ and $\sim 10^{-1}$ near threshold and near 100% SBS, respectively. It is contrary to our experimental observations / 13 / in which relatively large values of ion density modulation of about 10^{-1} near threshold have been measured. It is almost two orders of magnitude more than the prediction of Tajima model. Hence, this model is not applicable to our experiment as its predictions are not in good agreement with the experimental observations.

Evans / 14 / developed another model for the saturation of SBS. He assumed that the ion-waves generated by SBS may be saturated either by ion-trapping or by ion-Landau damping. His model is based on the assumption that the plasma has a constant electron temperature and a constant irradiance and SBS will increase with the scale length L . An increase in SBS will increase the ion temperature and ion waves will be severely Landau damped when ion temperature becomes equal to $T_e/2$. Hence, it will limit SBS to some maximum value:

$$R = \frac{2.4 \times 10^{16} \theta_o K_B T_e}{I_o \lambda_o^2} \quad (2.48)$$

where,

$$\theta_o = n_o / n_{cr}$$

$$I_o = \text{incident laser flux in } \text{Wm}^{-2}$$

$$\lambda_o = \text{laser wavelength in } \mu\text{m}$$

This model predicts $R=3.6$ for our parameters ($K_B T_e=52 \text{ eV}$, $\theta=0.16$, $I_o=5 \times 10^{14} \text{ Wm}^{-2}$) near threshold and $R=0.07$ for the maxi-

mum laser intensity available ($K_B T_e = 108 \text{ eV}$, $I_0 = 5 \times 10^{16} \text{ Wm}^{-2}$). Hence, this model also can not explain our results as it predicts a decrease in SBS with an increase in laser intensity, which is quite contrary to our experimental observations.

Kruer et al / 15 / presented a simple model of the laser plasma interaction in a frame moving with the reflection front. They showed that scattering of the incident light in plasmas can be limited by momentum and energy deposition as a result of the light. It does not describe a stationary state. Since a large fraction of the plasma will be pushed aside as the momentum and energy deposition from the damped ion waves secularly accumulates. This momentum deposition drives a reflection front supersonically through the underdense plasma whose speed increases as the square root of the laser intensity. Scattering depends on the laser intensity and on the density of plasma present in the underdense region. They predicted a decrease in the reflectivity as the reflection front moves through the plasma which in turn depends on the laser intensity.

This model is also not applicable as it predicts first a large reflectivity for a short time and then a decrease in the reflectivity with the increase in the laser intensity. Moreover, the backscattered light must have Doppler frequency shift. It is again contrary to our experimental observations {see section (4.2.2)}.

The above mentioned review of the theoretical models shows that no theoretical model exists which can explain our experimental results. However, an increase in SBS is expected on the basis of equation (2.46) both with the interaction length and with the density modulation which increases with the incident laser flux. In the derivation of equation (2.46) no saturation mechanism was taken into account. As the saturation mechanism can affect the backscattering to a large extent, they will be now briefly described.

2.2 Saturation Mechanisms

One of the most interesting problems in the field of LPI is the effectiveness of saturation mechanisms. Large amount of scattering does not imply a complete absence of all the saturation mechanisms.

It was found experimentally that the amount of the back-scattering does not increase so rapidly with the laser intensity as it is expected on the basis of the non-linear theory and it seems that at a certain stage saturation mechanism/s set in which affect the growth of SBS drastically. This can be easily inferred from SBS curves /16,17 / which show a sudden drop in SBS at a certain laser intensity.

2.2.2 Steepening of the Density Profile

Intense laser light impinging on a plasma may modify the density profile. This modification in turn affects the absorption, backscattering and the transport of energy. It is effective near the critical density /18,19/ as the amplitude of the es-wave may increase many fold due to the reduction of the group velocity from the velocity of light to the thermal velocity of electrons.

A lot of papers have dealt with /18-22/ the subject of profile steepening which is mainly theoretical and computational due to the stringent experimental requirements, for instance, spatial and temporal resolution of a few micrometers and a few picoseconds, respectively. Experimentally /21,22/ profile steepening has been observed near n_{cr} . It is not important at low densities, for instance below $0.4n_{cr}$ / 21 / and for short laser pulses. As we have deliberately used the underdense plasma ($n = 0.16 n_{cr}$) so this saturation mechanism can be excluded.

2.2.2 Rescattering

There is a possibility of the rescattering in the plasma which in turn can effect the growth of SBS. It can take place only when the backscattered wave amplitude in the plasma exceeds threshold for SBS. In such a case, the physical

mechanism responsible for rescattering is very simple. Suppose the pump wave is propagating in $\vec{k}_0$ direction and backscattering occurs somewhere inside the plasma. As the backscattered em-wave propagates through the plasma, it may excite an ion-acoustic wave in the $(-\vec{k}_0)$ -direction and a backscattered wave in the $(+\vec{k}_0)$ -direction. This process is known as rescattering. The rescattered wave in the forward direction (along $\vec{k}_0$) must have twice the frequency shift as the backscattered wave. This process may be repeated and is known as cascade process. Equations (2.40) and (2.41) can be modified to take into account the rescattering process. Its theory will not be described here as it has been discussed in detail by Speziale et al / 23 /. However, its presence will be checked experimentally in sections (4.2.2) and (4.2.3).

2.2.3 Secondary Decay of Ion-acoustic Wave

An ion-acoustic wave can decay into two further new ion-acoustic waves in an underdense plasma. The frequency of the ion-acoustic wave should be less than the characteristic frequency of the plasma. In our case, $\omega \ll \omega_{pi}$ hence such a decay is quite possible. It has been shown / 24 / that this secondary non-linear ion-wave decay is an efficient saturation mechanism for SBS in analogy with the lower hybrid wave heating / 25 /.

If the Vlasov-Maxwell equations are solved for the five waves (incident and backscattered em-waves, an ion-acoustic wave and its two daughter waves) then for an one-dimensional homogeneous plasma slab the threshold of ion wave decay by using the approximation of slowly varying amplitudes can be determined as / 26 /:

$$E_{th}^2 = \left(\frac{\Gamma_1 \Gamma_2 m_i c_s^2}{q_i^2} \right) \left(\frac{k^2}{k_1 k_2} \right) \quad (2.49)$$

where Γ_1, Γ_2 (k_1, k_2) are the phenomenological damping rates (wave numbers) of the daughter ion-acoustic waves, and E_{th} is the threshold electric field associated with the ion-acoustic wave. The threshold for the decay is minimum when $k_1 = k_2 = (1/2)k$. In our case, $E_{th} = 10^7$ V/m, while the threshold for SRS is $E_0 = 6 \times 10^8$ V/m.

So the threshold for the ion-acoustic decay is about two orders of magnitude smaller than SRS. On the other hand, the ratio of growth rates of SRS and secondary ion-wave decay is / 27 /:

$$\begin{aligned} \gamma_{IA} / \gamma_{SRS} &= \left(\frac{2 m_e \omega_o}{m_i \omega} \right)^{1/2} \left(\frac{\omega_o}{\omega_{pe}} \right) \frac{|E|}{|E_o|} \quad (2.50) \\ &= 6 \times 10^{-2} \quad (\text{near SRS threshold}) \\ &= 6 \times 10^{-3} \quad (\text{near 100\% SRS}) \end{aligned}$$

which shows that the growth rate for the ion-acoustic decay is much smaller than SRS. Hence, it is quite unlikely that in our case this phenomenon may occur.

2.2.4 Self Focusing and Filamentation

The laser focus is not necessarily a sharp point but can extend over a large region. Moreover, the focus may not be quite uniform, rather it may have several hot spots which may result in the filamentation instability.

Self-focusing of the incident laser beam in the subcritical plasma may occur. This is of concern in asymmetric fuel pellet implosion and the production of superthermal electrons that preheat the target. Evidences for laser beam filamentation have been reported in literature /28-29 /. It is easy to understand physically. In this process the ponderomotive force on the plasma by the laser radiation can cause the laser beam to break up into narrow light filaments which contain regions of locally much higher irradiance. Consider an initially incident laser beam (Fig. 2.6) that is slightly more intense in one region by an amount δI . The ponderomotive force $\vec{f}$ on the electron fluid due to this δI is to push the plasma away from the region of highest laser intensity (Fig. 2.6 b) leaving behind a region of lowered plasma density (Fig. 2.6 c). Since the index of refraction ζ of a plasma depends on the local electron density n_e as:

$$\zeta = 1 - \frac{n_e^2}{m_e \epsilon_o \omega_o^2} \quad (2.51)$$

the depletion of the electrons at the place where δI is large raises ζ there, which in turn reduces the phase velocity of the

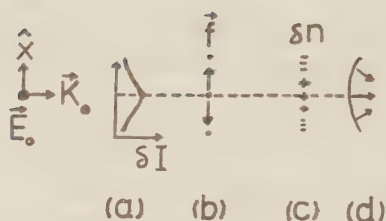


Fig.(2.6) Physical picture for the production of filamentation. A pondermotive force $\vec{f}$ due to laser intensity perturbation δI displaces the plasma away from the region of highest laser intensity, creating a channel of low density, and thus trapping the laser light and intensifies it. The enhanced beam can then drive more particles out of the region, and the process is repeated.

light wave and causes the wave front to curve as in Fig. (2.6 d). As a result, the light is focussed toward the original region of enhanced intensity, since the light wave energy propagates normal to the curved wave front. Thus an initial section of the beam with a slightly higher intensity than its surroundings will become more intense as the beam propagates. It is known as filamentation when it involves a small part of the incident light beam and is called self-focusing when it involves the beam as a whole.

As the characteristic dimension of self-focusing of the incident wave will be $\lambda_0 (c/c_s)$, so for an underdense homogeneous plasma, the threshold and the maximum growth rate of this instability is / 9 /:

$$\left(\frac{v_0}{c}\right)^2 > \left(\frac{w_0}{w_{pe}}\right)^2 \left(\frac{v_e}{c}\right)^2 \left(\frac{c_s}{c}\right)^2 \quad (2.52)$$

$$\gamma_{SF} = \left(\frac{v_0}{c}\right) w_{pi} \quad (2.53)$$

The threshold for self-focusing in our case is $\sim 10^{10} \text{ Wm}^{-2}$ and the inverse of the growth rate is about 200 ps, for a laser intensity of 10^{15} Wm^{-2} , which is about an order of magnitude larger than for SBS.

Liu and Ng /9,29/ believe that the filamentary structure is due to hot spots in the incident laser beam. We put a lot of effort to obtain an unstructured focal distribution. Baldis et al /30/ could confirm by an experiment that an unstructured focal distribution did not lead to self-focusing even up to flux values of 10^{17} Wm^{-2} in an underdense plasma.

In our experiment, self-focusing and filamentation can be easily ruled out on the following grounds:

- absence of hot spots in the laser output,
- within a short laser pulse of 2 ns, filamentation can not grow to a significant amount compared to SBS,
- our heating model /31/ shows excellent agreement between the cross-section of the interaction volume in the plasma and the laser focal spot diameter. In case of self-focusing, laser focal spot diameter will be much smaller and, hence, T_e , T_i will deviate considerably resulting in large discrepancies between theory and experimental results.

2.2.5 Ion Trapping

Nonlinear wave particle interaction leads to particle trapping. A particle moving in the direction of the propagation can be trapped and carried along by the wave until it is released by a collision.

All the ions moving with the velocity near the phase velocity of the ion-acoustic wave may be trapped in the wave potential. Hence, the condition for trapping is:

$$q \varphi_i = \frac{1}{2} m_i (c_s - v_i)^2 \approx \frac{1}{2} m_i \left(\frac{\omega}{k} \right)^2 \quad (2.54)$$

where, k , φ_i , and c_s are the wave number, potential and phase velocity of the ion-acoustic wave.

The bounce frequency ω_b of the ions, which is the frequency of the oscillations of the trapped ions at the bottom of the wave potential, can be determined as follows: The equation of motion of an ion in the wave is

$$\frac{\partial^2 x}{\partial t^2} = \frac{e}{m_i} E \sin kx \quad (2.55)$$

The ions are trapped in the negative part of the wave potential $\varphi_i = (E/k) \cos kx$ and a Taylor expansion gives:

$$\frac{\partial^2 x}{\partial t^2} + \omega_B^2 x = 0$$

where,

$$\omega_B^2 = \left(\frac{q k E}{m_i} \right) \quad (2.56)$$

The density perturbation in the plasma is (if $\delta n/n_0 \ll 1$)

$$\begin{aligned} \delta n/n_0 &= q\varphi_i / K_B T_e \\ &= q E / K_B T_e k \end{aligned} \quad (2.57)$$

Hence, the bounce frequency can be written in terms of the density modulations as:

$$\nu_B = \left(\frac{k}{2\pi} \right) v_e (m_e/m_i)^{1/2} (\delta n/n_0)^{1/2} \quad (2.58)$$

It shows that an increase in the density modulation or in the electron temperature will increase the bounce frequency of the ions. The density modulations have been measured by Gellert / 13 /, while the electron temperature can be determined by the following simple consideration:

The basic assumption of our heating model was that the plasma will be heated by laser light through inverse Bremsstrahlung mechanism. If Q is the local rate of heat flux absorption, then

$$Q = \alpha_h I_0 \quad (Jm^{-3}s^{-1}) \quad (2.59)$$

where, I_0 is the laser intensity in Wm^{-2} and α_h is the high frequency absorption coefficient / 32 /. So, (2.59) for $10.6 \mu m$ laser light can be written as:

$$Q = 9.7 \times 10^{-46} \ln \Lambda (K_B T_e)^{-3/2} I_0 n^2 \quad (2.60)$$

where, $K_B T_e$ is the electron temperature in eV, n is the plasma density in units of m^{-3} and $\ln \Lambda$ is the Coulomb logarithm :

$$\ln \Lambda = \ln \{ 12.3 \times (K_B T_e)^{3/2} \}$$

for $n = 1.6 \times 10^{24} m^{-3}$.

All the heat deposited by the inverse Bremsstrahlung will be transported away by the heat conduction as in a long cylinder. The heat equation is:

$$\frac{\partial}{\partial t} (1.5 n K_B T_e) + \nabla \cdot q = Q \quad (2.61)$$

where, q is the heat flux / 33 /

$$q = -3(n_e K_B T_e / m v_{ei}) \nabla K_B T_e \quad J m^{-2} s^{-1}$$

$$= -k_{th} \nabla K_B T_e \quad (2.62)$$

and k_{th} is the thermal conductivity coefficient

$$k_{th} = 3.1 \times 10^4 (K_B T_e)^{5/2} (\ln \Lambda)^{-1} \quad J s m^{-1} (eV)^{-1} \quad (2.63)$$

Hence, (2.63) can be written for the steady state case as:

$$-3 \times 10^4 (\ln \Lambda)^{-1} \{ (K_B T_e)^{5/2} \nabla (K_B T_e) \} = 9.7 \times 10^{-46} \ln (K_B T_e)^{-3/2} I_0 n^2 \quad (2.64)$$

The radial heat conduction equation will be:

$$(K_B T_e)^{3/2} r^{-1} \frac{\partial}{\partial r} \{ r (K_B T_e)^{5/2} \frac{\partial}{\partial r} (K_B T_e) \} = -2 \times 10^{-48} I_0 n^2 \quad (2.65)$$

where, $\ln \Lambda = 8$. Chen et al / 33 / have solved (2.65) by dividing the heated region in two zones. In the first zone ($r \leq b$), the laser intensity is assumed constant and finite, while in the second zone ($b < r < a$) laser intensity vanishes and over the distance of "a" T_e falls to the initial temperature of the plasma. It was sufficient to take $a/b=20$. The temperature at the beam boundary was estimated by comparing the heat flux at $r=b$ to the absorbed energy per unit length to be:

$$K_B T_e = 3.1 \times 10^{-10} n^{2/5} (I_0 A)^{1/5} \quad (2.66)$$

where, $A = \pi \varphi_0^2$ is the area of the cross section of the laser beam in the plasma and φ_0 is the laser focal spot radius in metre.

Equation (2.66) predicts an increase in the electron temperature with the laser intensity as shown in Fig.(2.7). It helps to estimate the density modulation and the bounce frequency of the ions.

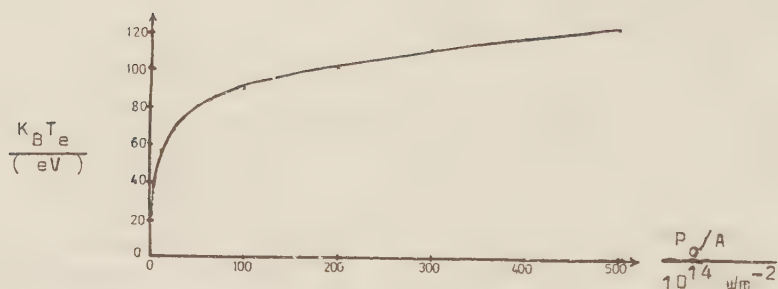


Fig.(2.7) Electron temperature as a function of laser intensity for a laser focal spot area of $7 \times 10^{-8} \text{ m}^2$ and a plasma density of $1.6 \times 10^{24} \text{ m}^{-3}$.

Electron thermal velocity (v_e) can be immediately determined from the electron temperature. The relationship is:

$$v_e = 4.19 \times 10^5 (K_B T_e)^{1/2} \text{ ms}^{-1} \quad (2.67)$$

where, $K_B T_e$ is in eV. For the conditions of interest 10-30% SBS, the corresponding laser intensity can be determined from the reflectivity curve (Fig.4.1). It is necessary to estimate the electron temperature from Fig.(2.7). It yields an electron temperature of about 75eV near 10% SBS which rises to 90eV for 30% SBS. The thermal electron velocity according to (2.67) corresponds to 3.6×10^6 and $3.9 \times 10^6 \text{ ms}^{-1}$ for 75 and 90 eV electron temperatures, respectively. As the electron thermal velocity does not change significantly in this range so an average electron thermal velocity of $3.8 \times 10^6 \text{ ms}^{-1}$ can be employed to determine the bounce frequency in the SBS range of 10-30%.

Besides the knowledge about electron thermal velocity, the amount of density modulation is also required for the estimation of the bounce frequency. It can be determined on the basis of the ion trapping model. Combining the equation (2.57) and the condition for the ion trapping (2.54) an analytic expression for the density modulation can be obtained:

$$\delta n/n_0 = \frac{1}{2} \left[\left(1 + \frac{3T_i}{T_e} \right)^{1/2} - \left(\frac{3T_i}{T_e} \right)^{1/2} \right]^2 \quad (2.68)$$

In Fig.(2.8), the density modulation has been plotted as a function of T_e/T_i . It shows that the density modulation strongly depends on the ratio between the electron and ion temperature. With the rise in the electron temperature the density modulation increases rapidly. It has been estimated / 31 / that in our case T_e/T_i will be in the range of 3-4 near 10% SBS. For this condition it can be easily concluded from Fig.(2.8) that the density modulation will be in the range of 8-10%. It is a little bit higher than the experimentally measured values / 13 /. A possible explanation will be given in section (4.4).

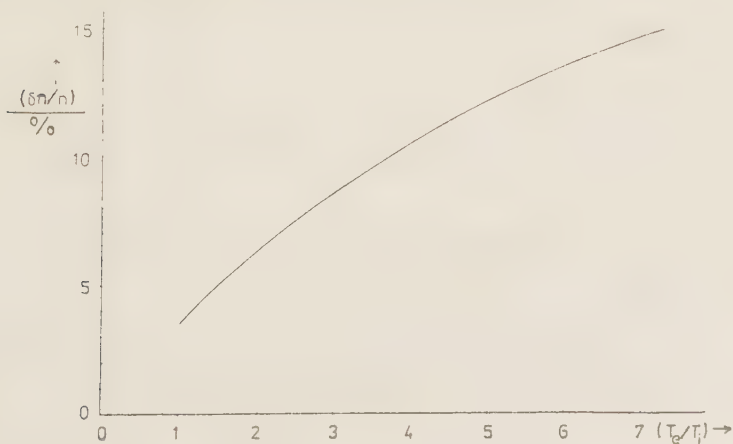


Fig.(2.8) Density modulation as a function of T_e/T_i as predicted by the ion trapping model.

By employing the average electron thermal velocity, the bounce frequency (2.58) can be written as:

$$v_B = 1.4 \times 10^{10} (\delta n / n_0)^{1/2} \quad (2.69)$$

An increase in the density modulation results in a gradual increase in the bounce frequency of the trapped ions, as shown in Fig.(2.9). It will be helpful in the interpretation of our experimental results (chapter 4).

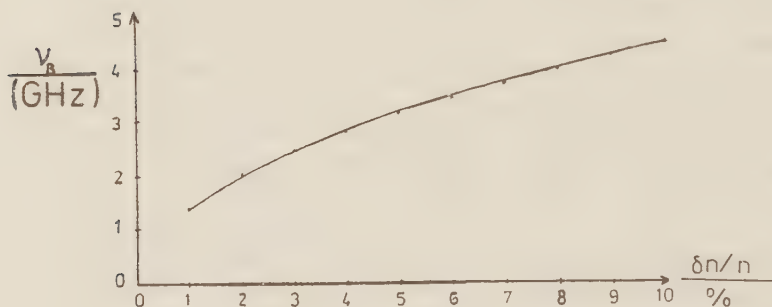


Fig. (2.9) Bounce frequency as a function of density modulation.

Now, the consequences of the ion trapping relevant to our case will be considered. They are:

One natural consequence of ion trapping is the energy gain by a fraction of ions. Ions with less energy will gain from the ion acoustic wave the energy (in the wave frame):

$$W_1 = \frac{1}{2} m_i (c_s - v_i)^2 \quad (2.70)$$

On the other hand, ions with more energy will give to the ion acoustic wave their energy:

$$W_2 = \frac{1}{2} m_i (c_s + v_i)^2 \quad (2.71)$$

The net loss of energy by the ion acoustic wave in its frame will be:

$$\Delta W = W_2 - W_1 = m_i c_s v_i \quad (2.72)$$

In our case, $v_i = 5.5 \times 10^4$ m/s, $c_s = 10^5$ m/s and the maximum velocity gained by the ions will be:

$$v_{im} = (2c_s v_i)^{1/2} = 10^5 \text{ m/s} \quad (2.73)$$

This shows that the particles can be accelerated up to the phase velocity of the wave.

When the ion acoustic wave amplitude grows to such an extent that the ion velocities are affected then the velocity distribution of the ions near the phase velocity of the ion wave changes. It will result in the appearance of a plateau in the ion Maxwellian distribution function, as shown in Fig.(2.10b). The main part of the ions velocity distribution function $f(v_i)$ remains unchanged. Only the ions in the vicinity of the ion acoustic phase velocity will be resonant with the wave.

As the trapping becomes significant, the particle orbits become non-linear. It is quite interesting to see how the motion of these particles evolves while the wave undergoes damped oscillations.

The behaviour of the ion trapping can be best understood in its frame by considering the ion orbits in the phase space, as shown in Fig.(2.10). The diagram shows the trapped particle orbits in the ion acoustic wave along with the relation between the wave potential ϕ_1 , trapping region in phase space and the ion distribution function. The ions in the region A are trapped but they may arrive at the point B (separatrix) by collisions. Then they will be accelerated ultimately to point C at the expense of the wave energy. They may now be decelerated to its original energy or may be untrapped by a collision. If the ions are deeply trapped in the wave potential, then the ion distribution $f(v_i)$ will

be large in region A, and the ion density of the trapped ions will be quite high at the bottom of the potential trough compared to marginally trapped ions near separatrix. The trapped ion distribution is shown in Fig.(2.10b). It shows a formation of the tail due to trapping in the ion-distribution function. The deformation of the velocity distribution is indicated in Fig.(2.10c). Ions in region 1 will be accelerated and those in region 2 will be decelerated resulting in a plateau in $f(v_i)$ as indicated by the dotted line. Computer simulations / 34 / have confirmed

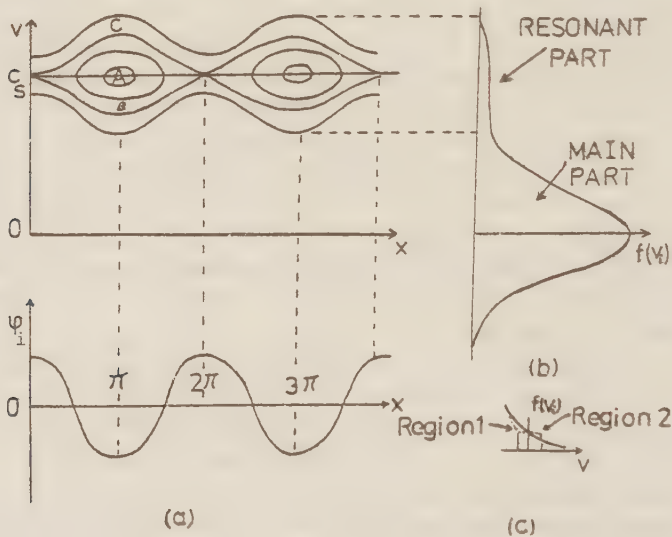


Fig.(2.10) (a) Trapped particle orbits in ion-acoustic wave along with the corresponding wave potential
 (b) division of the ion distribution into a main part and a resonant part, and
 (c) the deformation of the resonant part of the ion velocity distribution $f(v_i)$.

that the particles are accelerated by the trapping process and a fraction of the ions are heated forming a tail in the ion velocity distribution function. This deformation of $f(v_i)$ is quite significant as it affects the growth of the ion acoustic wave and will be discussed later.

Ions will be accelerated at the expense of the ion-acoustic wave energy. It can be roughly estimated as: From (2.72), the loss of energy by the ion acoustic wave can be estimated for our parameters as:

$$\Delta W = 9.4 \times 10^{-18} \text{ J} \quad (2.74)$$

For a density modulation of 4% and for a particle density of $1.6 \times 10^{24} \text{ m}^{-3}$, it corresponds to

$$\delta n \Delta W = 4.4 \times 10^5 \text{ J m}^{-3} \quad (2.75)$$

The laser focal spot area was about $7 \times 10^{-8} \text{ m}^2$ and if the interaction length is 1mm, then the total power lost by the ion acoustic wave within a half bounce period of 100 ps is:

$$P_1 = 3 \times 10^5 \text{ W} \quad (2.76)$$

On the other hand, the laser power (say 1GW in a pulse) conversion into the ion acoustic wave can be determined from the Manley-Rowe relation as:

$$1 - \frac{(w_o - w)}{w_o} = 0.00056$$

Hence, the power of the ion acoustic wave will be

$$P_m = 5.6 \times 10^5 \text{ W} \quad (2.77)$$

From (2.76) and (2.77), it can be concluded that trapped ions can take about 50% of the wave energy resulting in the damping of the wave.

Another consequence of the ion trapping may be an additional frequency shift in the backscattered light as has

been calculated by Shukla /35/ who extended the work of Schamel /36/. The physical basis of the shift in the frequency of an ion acoustic (IA) wave is that the potential of the IA-wave becomes sufficiently large in order to reflect the particles, then consequently the frequency of the wave becomes a function of amplitude in the nonlinear regime. This effect originates because the distribution function and hence the number density does not remain linear during the trapping process. In case of large amplitude waves, the dispersion relation for the nonlinear case should be solved. It gives an additional frequency shift ($\delta\omega$) due to trapped ions as:

$$\delta\omega = -\frac{16}{5} (\delta n/n)^{1/2} (k\lambda_i)^{-2} \left(-\frac{T_i}{T_e}\right) a_e \left[1 + \frac{3}{2} a_i\right] (\partial\epsilon / \partial\omega)^{-1} \quad (2.78)$$

where,

$$\partial\epsilon / \partial\omega = 0.5(T_e/T_i) \frac{(k\lambda_{de})^{-2}}{\omega} \left[(1+k^2\lambda_{de}^2)(T_i/2T_e)(4\xi^2-1)^{-1} \right] \quad (2.79)$$

$$\xi = \frac{\omega}{2kv_i}$$

$$a_e = \frac{4}{3\pi} (\omega/kv_e)^2 \left[\exp(-0.5(\omega/kv_e)) \right] \quad (2.80)$$

$$a_i = \frac{4}{3\pi} (T_e/T_i)^{1/2} (\omega/kv_i)^2 \left[\exp(-0.5(\omega/kv_i)) \right]^2$$

In our case, the additional non-linear frequency shift is about 7% for a density modulation of about 8%. It is quite small and can be neglected in our case as for a frequency shift of (17 ± 1) GHz (section 4.2.2). It corresponds to ~ 1 GHz which is within the experimental error.

Before, the ion trapping model is applied to our case it is necessary to check that the threshold criterion is satisfied. So, it will be described in the next section.

2.2.5.1 Threshold of Ion Trapping

Landau / 37 / has discussed the case of a collisionless damping of linear plasma oscillations. He considered only the straight line particle orbits and neglected the particle trapping effects. As the amplitude of the ion acoustic wave increases, it may trap some ions resulting in the deformation of the Maxwellian distribution of the ions (Fig.2.10). The zero order velocity distribution is flattened in the vicinity of the phase velocity of the ion-acoustic wave. According to linear theory, Landau damping should disappear which is contrary to the experimental observations / 38 - 41 /. Hence, the trapping process has to be treated by a non-linear theory.

A lot of theoretical work has been done to see the effect of particle trapping on the electron/ion wave in the plasma. There are two major approaches. In the first approach, as followed by Knorr / 42 /, and Armstrong / 43 /, Vlasov and Poisson's equations were solved numerically. In the second approach, as followed by O'Neil / 44 / and Dawson /45,46/, the electron distribution function was divided into resonant and non-resonant parts and then the equation of motion was solved exactly for the trapped electrons moving in their self consistent field. The two methods seek the same ends.

First of all, O'Neil considered the case of electron trapping in the potential well of the wave. He found that after some initial Landau damping the wave amplitude is being modulated with the bounce frequency. These modulations reach a constant value after a long observation time. However, in his treatment, the ion-acoustic wave amplitude was assumed to be constant. This approximation does not affect the results seriously for large amplitude waves. An exact treatment has been performed by Oei et al / 47 / and Franklin et al / 41 /. They consider the case of a wave whose amplitude is slowly time varying. Oei et al established the threshold condition for the electron trapping by non-linear plasma waves as $g_c \leq 0.6$, where g_c is the critical value of g and it is the ratio between the Landau damping rate and the bounce frequency. They found that the small amplitude waves damp exponentially ($g \geq 0.6$) and large am-

plitude waves ($g \leq 0.6$) undergo amplitude modulations. Their results were confirmed by the computer simulation of Sugihara et al / 34 /. They found the critical value of g as 0.77. The results of Sugihara et al extend over much longer time than in earlier work. Below the critical value of g the amplitude of the wave becomes constant or damps after several modulation periods. Similar results were obtained by Canosa et al / 48 /. It was also confirmed experimentally /38-41, 49 /. Such oscillatory behaviour results in the appearance of the side bands which will be discussed in the next section.

The value of g depends directly on the density modulation. In Fig. (2.11) g is plotted as a function of the density modulation for our case. Here, g is the ratio between the total damping rate and ω_g . As the critical value of g is 0.77 / 34 /, so, in our case, the minimum density modulation required for the ion trapping to be effective is about 8%. It implies that the ion-acoustic wave amplitude will oscillate if the density modulation is more than 8%.

In addition to the amplitude modulation of the ion-acoustic wave on the long time scale, the enhanced initial damping was predicted by Dawson et al / 46 / and Armstrong / 43 /. They found that the damping of a large amplitude wave is much stronger than that predicted by the linear theory. This enhanced damping lasts only while the particles are being accelerated, and then saturation ultimately occurs. Such an enhanced damping in case of an ion acoustic wave was observed by Sato et al / 38 / for the initial density perturbation of more than 10%. They also observed the decrease of the damping after this strong damping and the amplitude modulation of the larger amplitude waves. Moreover, they observed that even a small amplitude wave shows modulations in the presence of the large amplitude wave. They attributed it to the deformation of the velocity distribution by the large amplitude wave.

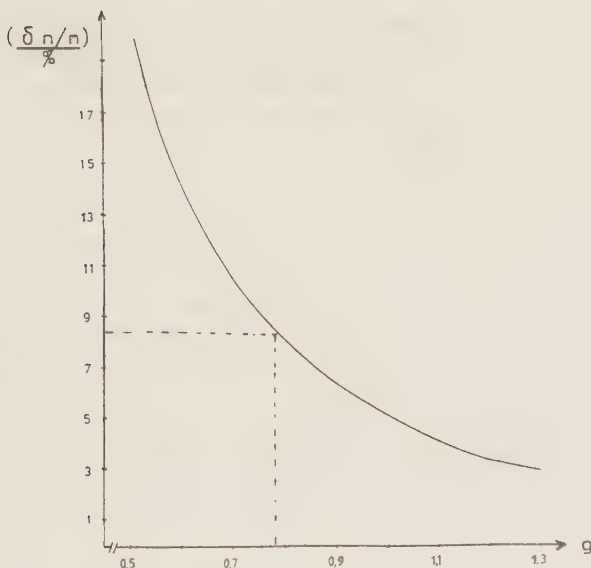


Fig.(2.11) g is plotted as a function of the density modulation $\delta n/n_0$, where g is a ratio between the damping rate and the bounce frequency. The critical value of g is 0.77 above which particle trapping is not effective.

Besides enhanced damping the ion acoustic wave is modulated due to ion trapping which will result in the generation of the side bands to be discussed in the next section.

2.2.5.2 Side Bands

It is now well established /39, 41, 45, 50-52/ that a finite amplitude monochromatic electron wave propagating in an one dimensional plasma may generate side bands. These side bands result from the motion of the trapped electrons in an es-wave. First of all, Wharton et al / 39 / observed the side bands to the frequency of the main wave in case of electron trapping.

The frequency separation was found to be proportional to the square root of the density modulation. This indicates that the trapped particles bouncing in the potential trough of the wave at a frequency of ω_B play an important role in the generation of the sidebands.

Physically, the interaction between a modulated wave and the trapped particles bouncing in its potential trough will result in sidebands. The sidebands extract the energy from the large amplitude wave and cause it to be heavily damped.

On the other hand, amplitude modulations in the ion-acoustic wave while propagating through the plasma have been confirmed experimentally /38,49,53-55/. Similar to the case of electron trapping, the sidebands in case of ion trapping have been observed by Ikezi et al /53/. They observed the distinct sidebands when $\delta n/n_0$ exceeds 2%. Moreover, they also found an enhanced damping when the sidebands exist than when it is absent.

The ratio between the amplitude of the upper and lower sidebands can be well explained by the trapping particle model. A general theory of the trapped particle instability has been developed by Mima and Nishikawa /56/. They found that both the upper and lower sideband amplitudes depend on ΔN_t , where $\Delta N_t (=N_a - N_0)$ is the "excess" trapped particle density, N_a is the actual number density of the particles and N_0 is the particle density according to the Maxwellian distribution. The upper (lower) sideband has greater amplitude if $\Delta N_t > 0$ ($\Delta N_t < 0$). It gives some insight into the ion wave dynamics, for instance, in the damping of the wave which in turn depends on the number of the particles trapped in its potential. This criterion will be employed to explain our experimental results in section (4.4).

2.2.5.3 Effect of Collisions

It must be pointed out here that all the above analysis assumed a collisionless plasma which was not the case in our experiment.

The effect of collisions leads to a smoothening of the distribution function. Mima and Nishikawa /56 / considered the effect of collisions in the case of electron trapping. If it is extended to the case of ion trapping by including the ion dynamics the effect of collisions will become significant when

$$\frac{v_{ii}}{w_B} > \left(\frac{c_s}{v_i} \right)^2 \left(\frac{w_B}{w} \right)^2 \quad (2.81)$$

where, v_{ii} is the ion-ion collision frequency and v_i is the ion thermal velocity. The condition can be written as:

$$H = \left(\frac{c_s}{v_i} \right)^2 \left(\frac{w_B}{w} \right)^2 \frac{w_B}{v_{ii}} < 1 \quad (2.82)$$

It implies that the collisions will become ineffective as soon as H becomes greater than unity. H is plotted versus density modulation for our parameters in Fig.(2.12). It shows that collisions will be effective for a density modulation of less than 9%.

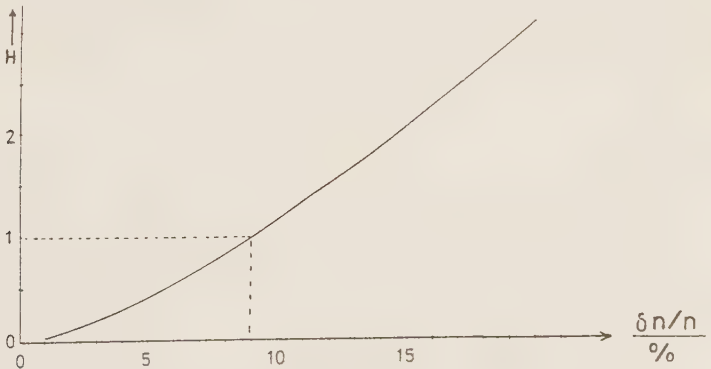


Fig.(2.12) H is plotted as a function of the density modulation. The dashed line shows the limit below which collisions will affect the ion trapping.

2.2.6 Ion Acoustic Wave Harmonics

A large amplitude ion acoustic wave will be produced in the plasma during a laser plasma interaction experiment. It will be a monochromatic pure sine wave as shown in Fig.(2.13a). Ion trapping will set in as soon as this wave will become sufficiently large or in other words when the threshold condition for ion trapping is satisfied, as mentioned in section (2.2.5.1). The wave will influence the motion of the particles. At the first stage, some particles will slip out of the sinusoidal wave shape as shown by a dotted curve in Fig.(2.13b). It will result in the production of the harmonics.

Karttunen and Salomaa / 57 / found that the generation of ion wave harmonics is a loss mechanism resulting in the damping of SBS. Their analysis is based on the assumption of a perfect matching between the ion wave and its second harmonic, while neglecting the higher order harmonics. Moreover, no damping term for the second harmonic was considered. His simple model predicts a gradual increase in the back-scattering at low plasma density ($n < 0.2n_{cr}$) with the laser intensity. It will be further discussed in section (4.4).

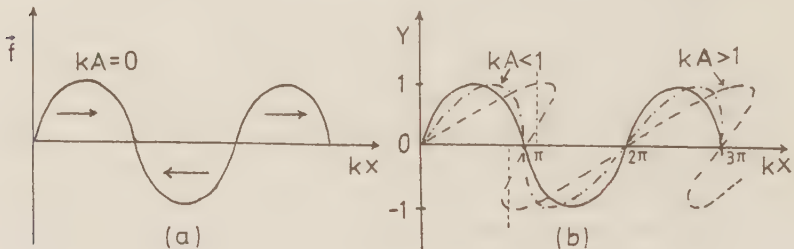


Fig.(2.13) (a) Force on the plasma particles forcing them to change their velocities.
 (b) Sequence of events leading to wave breaking. Initially the ion acoustic wave is sinusoidal ($KA=0$) and then it distorts ($KA < 1$) due to ion trapping leading to even wave breaking ($KA > 1$) at later stages.

2.2.7 Wave Breaking

If the amplitude of the ion acoustic wave is very large so that the order of particles along k is too much disturbed (Fig.2.13b) then wave breaking will occur. The condition for the wave breaking can be determined by the modification of the Hughes model / 58 /.

Consider motion only in the x -direction and assume that the ions oscillate in the ion acoustic wave field with a displacement $X(x_0)$, where x_0 is its mean position. The displacement of ions will result in an electric field:

$$E = - \frac{n e q}{\epsilon_0} X \quad (2.83)$$

The equation of motion for ions, under the assumption of parabolic wave potential, will be:

$$\frac{\partial^2 X}{\partial t^2} + \omega_{pi}^2 X = 0 \quad (2.84)$$

where, ω_{pi} is the ion plasma frequency. This model is valid for the small displacement (X) of ions, where the particle position is given by:

$$x = x_0 + X \quad (2.85)$$

If the ions are trapped in the ion acoustic wave then in the stationary case:

$$X = A \sin kx_0 \quad (2.86)$$

where, A is the maximum displacement during bouncing. By defining $Y = E/E_0$ and employing (2.83) equation (2.86) can be written as:

$$kX = (kA)Y + \sin^{-1} Y \quad (2.87)$$

Y is plotted versus kx in Fig.(2.13b). It shows that, if $kA > 1$, then Y becomes double valued near $kx = (2n+1)\pi$. This is physically not possible, hence wave breaking occurs. For no wave breaking:

$$kA \leq 1 \quad \text{or} \quad \omega_B^2 \leq \omega_{pi}^2 \quad (2.88)$$

In our case, ω_B is orders of magnitude smaller than the ion plasma frequency. So, there is no chance for wave breaking.

2.2.8 Ion Heating

Ion heating is due to the resonant interaction of an ion acoustic wave with the ions which result in the uniform heating of the bulk ions rather than a fraction as it is in the case of ion trapping. Here, a simple model of Krueer/10 / will be considered, which takes into account the finite amount of ion heating.

In the case of ion heating, first of all resonant ions will gain energy from the ion acoustic wave which will be further transfer to the nonresonant ions through collisions. This process will continue untill the ion acoustic wave is fully damped and the bulk of ions are uniformly heated. In this way, it may be a very efficient saturation mechanism.

To find out the amount of backscattering in the presence of ion heating as a saturation mechanism, ion temperature should be estimated. To predict the ion temperature, it is assumed that all of the ion acoustic wave energy is going into heating the ions and that the energy is being transported away out of the interaction volume by simple heat conduction mechanism. For such a case, Chen / 59 / found an equation for the ion temperature as:

$$(K_B T_i)^2 (K_B T_i - K_B T_0) = 3.6 \times 10^{-9} \left(\frac{\varphi_0}{L} \right) R I_0 \quad (2.89)$$

where, φ_0 is the radius of the laser focal spot, $K_B T_0$ is the initial ion temperature, L is the interaction length and R is the amount of backscattering.

The above equation was solved for an initial temperature of 10eV and for a focal spot radius of 150 μm . The results are shown in Fig.(2.14) which shows that the ion temperature rises sharply with the laser intensity. It also shows a dependence of T_i on the interaction length.

The size of the ion acoustic wave can be calculated from (2.22). Keeping only the damping term on the left hand side and the resonant terms on the right hand side, it results in:

$$\delta n/n_0 = \frac{e^2 k^2}{2m_e m_i \Gamma_i \omega} A_0 A_- \quad (2.90)$$

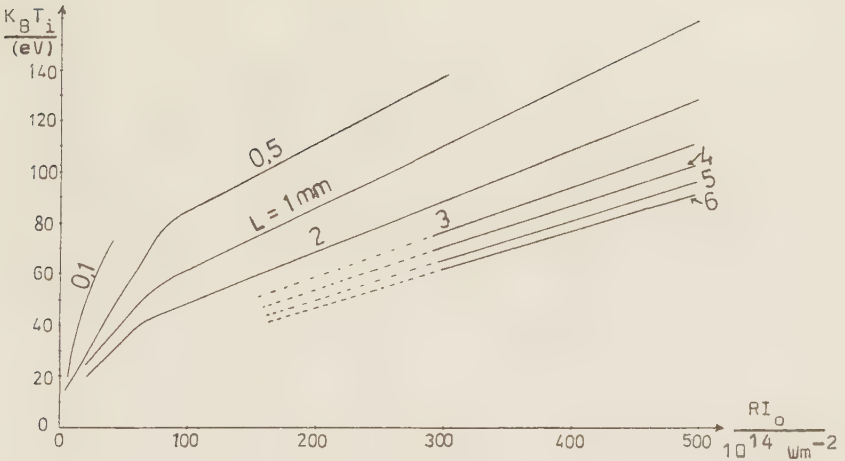


Fig.(2.14) Ion temperature as a function of RI_0 as predicted by the ion heating model.

where, Γ_i is the Landau damping decrement for the ion acoustic wave and ω is its frequency. If the above mentioned equation is solved along with the equations (2.40) and (2.41), then the amount of backscattering can be determined from / 10 /:

$$R (1 - R) = B [\exp\{G'(1-R)-R\}] \quad (2.91)$$

where,

$$G' = 4.76 \times 10^{-2} \left(\frac{k_o L}{\Gamma_i} \right) \omega \eta^2 \left(1 + \frac{3T_i}{T_e} \right)^{-1} \quad (2.92)$$

for our parameters. Here, B is the initial noise level, k_o is the wave number of the incident wave.

The knowledge of T_e and T_i from (2.66) and (2.89) will enable to calculate the amount of the backscattering for the case of ion heating as a possible saturation mechanism.

The efficient ion heating reduces T_e/T_i and consequently Landau damping becomes effective. Therefore, the model predicts a low amount of backscattering. An increase in the backscattering with the scale length is much slower than was predicted by (2.46) when ion heating was ignored. Chen / 59 / made an attempt to explain his and our experimental results on the

basis of this model, but did not succeed. Near threshold, the interaction length will be small resulting in a sharp increase in T_i { Fig.(2.14) } with the laser intensity. It is contrary to our experimental observations (section 4.2.2), so ion heating can not be a possible saturation mechanism in our experiment causing anomalous damping near 10% SBS.

Now, there is one more possibility which could affect SBS and will be discussed in next section.

2.2.9 Pump Depletion

As the amplitude of the backscattered wave grows and becomes equal to the incident laser wave amplitude, the pump wave will be depleted.

Fuchs presented an analytical study of SBS including the effects of pump depletion. His model / 60 / predicts an increase in the backscattering rapidly towards unity for a constant damping rate (Γ) with the normalized plasma length which in turn depends on the laser intensity. Backscattering seems to saturate at an early stage as Γ increases.

Forslund / 8 / also considered the effect of pump depletion as the amplitude of the backscattered wave increases. His computer simulation results show a decrease in the amplitude of the incident pump wave and a frequency modulation in the backscattered wave after it has reached the saturation stage. Pump depletion will be effective only when SBS is near 100% and hence can not effectively reduce SBS, rather it may result in frequency modulations. Such frequency modulations were predicted on a time scale of a few picoseconds which are beyond the resolution of our IRSA system.

Chapter 3

EXPERIMENTAL SET-UP

The aim was to investigate the scattering processes, in particular SBS, under well known laser and plasma conditions. For this purpose, a high power CO₂ laser system was developed to investigate a preformed Z-pinch plasma. In the following sections, the high power laser system, the Z-pinch, and the procedures and equipments used for the scattering experiment will be described.

3.1 The High Power CO₂ Laser System

A high power CO₂ laser system was developed here to study SBS resulting from laser plasma interaction. The laser output must possess certain basic properties for a clean LPI experiment. For instance, high contrast ratio, single longitudinal and transverse mode, small beam divergence and variable pulse-width. Moreover, the laser pulsewidth should not be greater than 2ns to avoid any density variation in the plasma. The ion temperature was 16eV (section 4.2.2) near threshold of SBS corresponding to the ion thermal velocity of $5.5 \times 10^4 \text{ ms}^{-1}$. Hence, the ions could escape from the interaction volume of 200 μm diameter (page 54) within 2ns.

Sharp laser pulses were achieved by injection mode locking (IML) as described in the appendix. It is based on the principle of injecting a pulsed signal from a laser termed as master oscillator into another laser oscillator (slave oscillator). As a master oscillator, a helical TEA CO₂ laser of the type investigated in / 61 / was employed.

Fig.(3.1) shows the essential features of the helical TEA laser. The discharge vessel consisted of a 3 metre long and 37 mm diameter plexi-glass tube. A large number of 1 K Ω and 10 Ω resistors were glued in the discharge tube by epoxy to preclude arc formation. Since a glow discharge at atmospheric pressure is unstable and rapidly develops into an arc. As shown in Fig.(3.1a), the discharge current was fed to each point electrode through a resistor. These resistors tend to

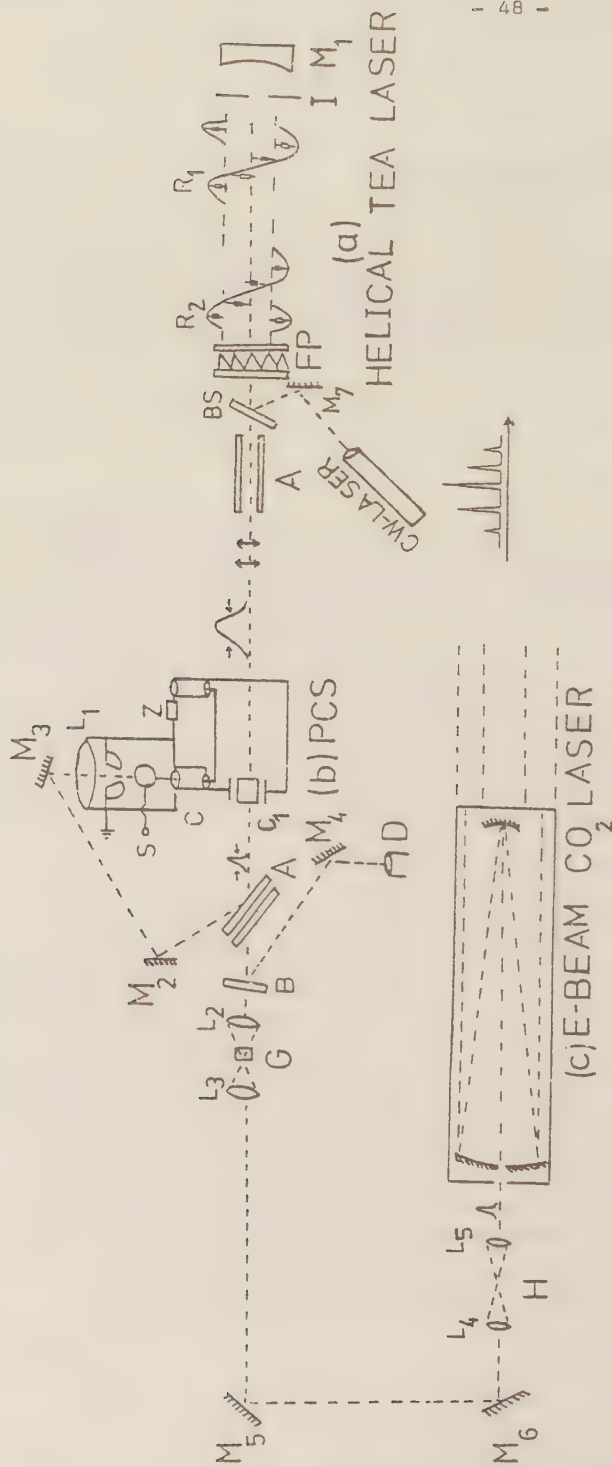


Fig. (3.1) The high power CO₂ laser system consisting of different lasers; a pulsed helical TEA laser, a CW laser and an e-beam laser. Here, M's are gold coated mirrors, I is an iris, R₁ (R₂) are 10Ω 1(KΩ) resistors, FP is a thermal resonance reflector, BS is a beam splitter, A is a polarization analyser, C₁ is a CdTe crystal, C is a polarity inverting cable, S is the high voltage terminal of the spark gap, Z is a 50Ω cable termination, L's are the lenses, G is a SF₆ cell, D is a detector, B is a ZnSe window, and H is an optical isolation system.

stabilize the discharge by preventing large currents from being drawn through any one pin electrode. A 5.6 m radius of curvature almost totally reflecting (98.5% reflectivity at $10.6 \mu\text{m}$) gold coated mirror was employed on the rear of the laser resonator. It oscillates at a number of longitudinal modes. A typical uncontrolled output is shown in Fig.(3.2a). The output was modulated due to the beating between different longitudinal modes. It was highly undesirable, as the aim was to inject a single frequency pulse into an e-beam laser. The single line operation was achieved by employing a resonance reflector as a front mirror. It is made of Germanium flats of 10 mm thickness and was mounted in a holder in which water was circulating and whose temperature was kept constant by a thermostat. Laser operation at different lines can be obtained by simply tuning this Fabry Perot. Moreover, the helical TEA laser was forced to oscillate at TEM_{00} mode by introducing an iris inside the resonator thus selecting an axial mode by diffraction losses.

The laser cavity has Brewster windows. It yields modulated output polarized in the vertical direction with the output energy of a few tens of milijoules. It has a pulsewidth of about 200 ns and a beam divergence of 1.5 mrad.

A pulse cutting system (PCS), as described in / 61 /, was employed to obtain a short pulse of a few nanosecond duration for IML. This system consisted of a laser triggered spark gap, a CdTe electrooptic crystal and a polarization analyzer, as shown in Fig.(3.1b). The high voltage electrode S was connected over a short cable to the CdTe crystal. The radiation reflected from the analyzer A was directed to the spark gap. A 50 mm Ge lens was employed to focus the radiation to the electrode S. The breakdown of the voltage occurs with subnanosecond risetime applying a high voltage at the CdTe crystal. This rotates the polarization direction of the light for a very short time resulting in the inverse of a short pulse to the transmission direction of the analyzer. The pulsewidth could be varied by changing the cable length of the spark gap.

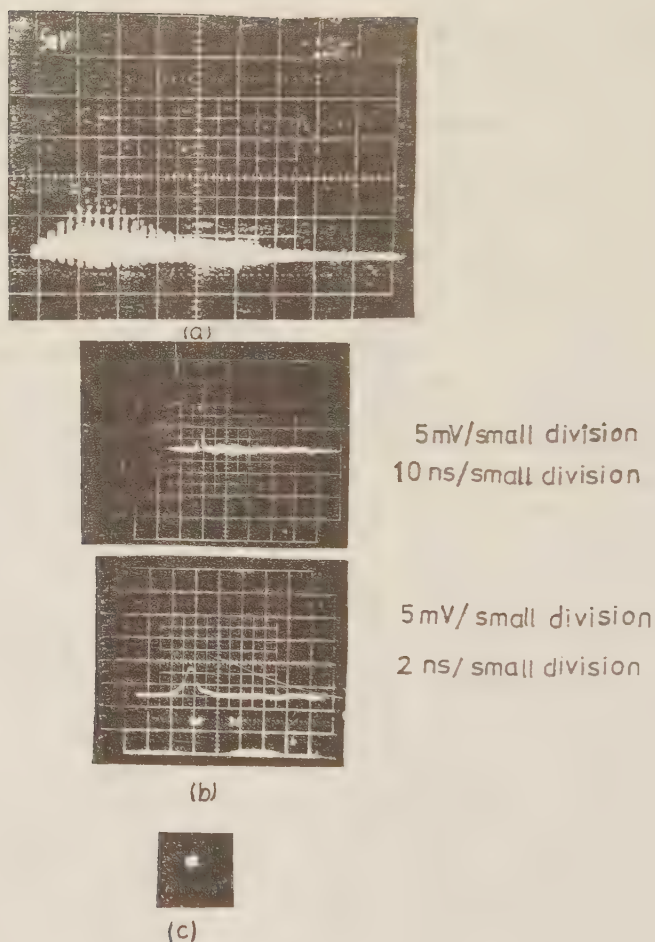
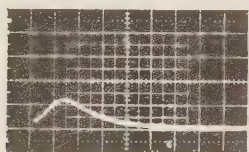


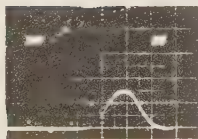
Fig.(3.2) (a) Modulated output of the helical TEA CO₂ laser,
(b) a sharp output pulse delivered by the pulse cutting system on two different time scales, and
(c) the near field distribution.

For the shortest cable length, a 2-3 ns short pulse was obtained as shown in Fig.(3.2b).

By increasing the cable length of the spark gap of the PCS, a laser pulse longer than 2 ns was obtained. But it was spiky in nature. To obtain long laser pulse without modulations from the helical TEA laser a CW laser (see specifications in appendix) was coupled into it by a salt window (Fig.3.1). It does not change the setting of the system when the signal was fed in. In this case, the emission of the TEA laser was a smooth pulse of about 200 ns as shown in Fig.(3.3a). The PCS reduces it to a few nanoseconds, Fig.(3.3b).



(a) 200 Kw /
100 ns / small division



(b) 20 Kw /
1 ns / small division

Fig.(3.3) (a) Helical TEA CO₂ laser output with mode locking
(b) the same output but shortened by PCS.

On the other hand, the slave oscillator is an e-beam Maxwell CO_2 laser system which consisted of an electron beam gun, a discharge chamber, a Marx generator, and a sustainer bank. It is schematically shown in Fig.(3.4).

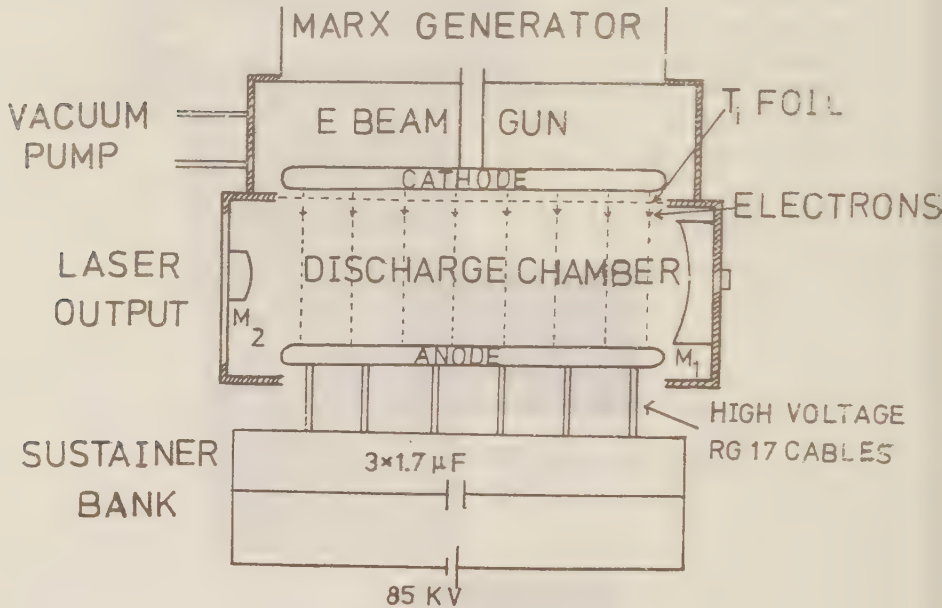


Fig.(3.4) A Schematic diagram of the e-beam laser system.

The cold cathode chamber consisted of a stainless steel vacuum vessel (0.61 m diameter, 1.52 m length) and a large cathode capable of withstanding negative high voltage of up to 300 KV for 10 μs . The vessel is designed for an operational vacuum down to 10^{-3} Pa. It has a cold cathode electron gun. In contrast to the thermionic cathode / 62 /, which is capable of supplying moderate current densities ($\leq 10^4 \text{ Am}^{-2}$) for arbitrary pulse lengths, the cold cathode operates satisfact-

orily at high current densities ($\geq 10^4 \text{ Am}^{-2}$) and for times less than a few μs . The cold cathode gun is simple to construct and easy to operate compared to thermionic cathode, requires no auxiliary power supplies or very high vacuum technology, and can not be damaged by vacuum failure.

The entire e-beam high voltage is housed in an oil filled steel tank and the source is a 200-300 KV Marx generator. It consists of four stages with a maximum stage voltage of 75 KV and a stage capacitance of $1.2 \mu\text{F}$. Each stage comprises of a three electrode spark gap switch and charging and triggering resistors. A Marx bank was used instead of a transformer as it has a low inductance, small energy leakage and is more efficient.

Parallel to the cathode is the window through which electrons enter in the gas chamber. The gas chamber and the e-beam system are separated by a thin foil, which can withstand a high pressure. The window is a 200 μm to 40 μm thick titanium foil.

The discharge (or gas) chamber is a pressure vessel made of mild steel having a diameter of 0.45 m and a length of 0.164 m, one large area carbon anode of $0.01 \times 0.11 \text{ m}^2$ and one adapter plate which mates the gun chamber. It can withstand a pressure up to about 10^6 Pa .

The sustainer bank required to sustain the gas discharge initiated by the 200-300 KeV electron pulse consisted of three 60 KV capacitors of $1.7 \mu\text{F}$ each which were charged by a power supply delivering up to 85 KV at 10 mA. The capacitor bank was connected directly to the discharge electrodes through six coaxial cables in parallel. Because the large cross-section laser could bear high emission losses, the unstable resonator configuration depicted in Fig.(3.4) was chosen. It was made up of the concave copper mirror M_1 and a convex one M_2 of 10 and 5 m radius of curvature, respectively, separated by 2.5 metre and with their centres of curvature coinciding it, thus, was of the so called confocal type and with the diameters of 100 and 50 mm, respectively, had an output coupling of 75%. The laser output leaving the cavity as a plane wave outside the convex mirror through a salt window showed in this way an annular field distribution, Fig.(3.5).

The discharge chamber of the e-beam laser was filled with a gaseous mixture of He and CO_2 at a pressure of 2×10^5 Pa. The gaseous mixture resides in a steady electric field E_s of the sustainer bank and is irradiated by a high energy pulsed electron beam. Each primary electron creates a swarm of secondary electrons in the gas generating a uniform discharge in a quite large active volume of 10 litres.

The near field intensity distribution was observed by burning the polaroid film, as shown in Fig.(3.5), while the far field intensity distribution was recorded by the IRSA-system and also by burning a polaroid film as shown in Fig.(3.6) and is described in detail in / 31 /. It shows that the focal spot radius of the laser was about $50 \mu\text{m}$ in the focal plane of a 714 mm lens. A 1500 mm lens was sometimes employed to focus the laser radiation. In this case, the radius of the focal spot was found to be about $100 \mu\text{m}$. A lot of effort was invested to avoid any substructure in the laser focal spot (section 2.2.4). It was achieved by adjusting the laser system to deliver radiation which was diffraction limited by the output window of the unstable resonator resulting in a bell shaped focal distribution of the above dimension.

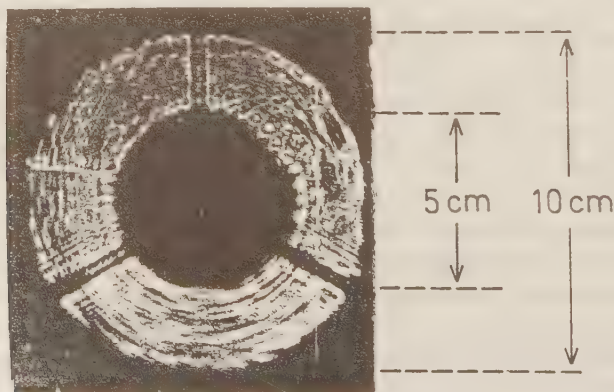


Fig.(3.5) Annular field distribution of the e-beam laser.

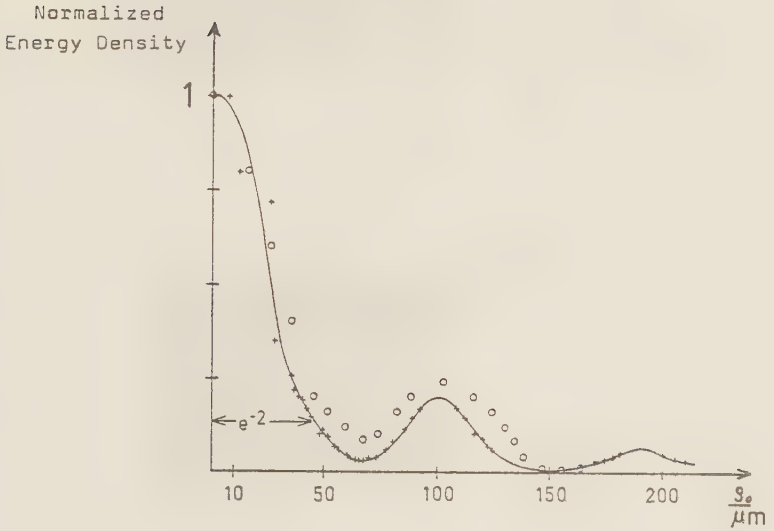
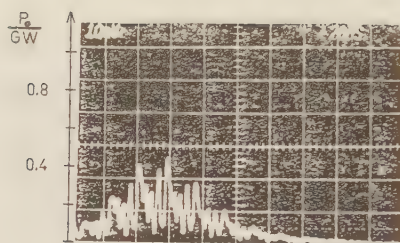
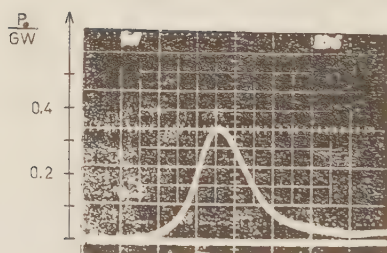


Fig.(3.6) Far field energy distribution of the e-beam laser. Circles and crosses represent the IRSA-system and the polaroid film burning pattern results, respectively. The e^{-2} focal spot diameter is about $100 \mu m$ in the focal plane of a 714 mm lens and has no substructure.

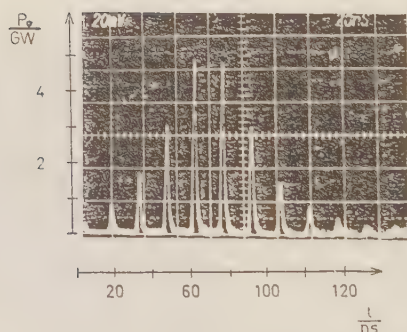
In the absence of injection mode locking, the output power was found to be about 500 MW with a pulsewidth of about 70 ns as shown in Fig.(3.7a). Mode beating results in an modulated output. The beam divergence was found to be 0.06 mrad.



(a)



(b)



(c)

Fig.(3.7) E-beam laser output; (a) uncontrolled output, (b) when a cw CO_2 laser was coupled into the e-beam laser (injection locking), (c) when a short laser pulse was coupled in (injection mode locking). p_o is the laser output power.

For IML, a few nanosecond pulse from the master oscillator was injected into the slave resonator through a hole in the centre of the rear mirror of the e-beam laser. The e-beam laser then emits a train of nanosecond pulses with an inter-separation time of 16.7 ns, which is equal to the round trip time of the resonator, Fig.(3.7c). The complete high power CO₂ laser system is shown in Fig.(3.1). An SF₆ cell was employed between the master and the slave oscillator to increase the contrast ratio and to avoid prelasing coupling between master and slave laser. An optical isolation system, which consisted of two lenses, was used to protect the SF₆ cell and other components from the high power laser radiation being emitted backward from the coupling hole in the rear mirror of the e-beam laser.

For the experiment to be performed, the laser output power was important rather than the energy. So, a gas mixture of He and CO₂ was used and N₂ gas was deliberately avoided in order to reduce the energy of the delayed pulses which may damage optical components. Sometimes, a CW CO₂ laser output was directly coupled into the e-beam laser. The output was then found to be an unmodulated pulse of about 70 ns, as shown in Fig.(3.7b). Threshold for injection locking was determined by attenuating the injected signal. It was found to be in the range of 0.4 - 2.5 Wm⁻². Injection locking occurs in 10% cases when the injected signal power density was 0.4 Wm⁻² and in 90% cases when it was 2.5 Wm⁻².

The output of the helical TEA laser was recorded by a pyroelectric detector P5-01 with sufficient time resolution and the e-beam laser output was measured by the subnanosecond photon drag detectors PD-7410 and PD-7455. Spectral measurements were made by a grating spectrometer equipped with an time integrating so called thermal image plate in the output plane (Opt. Engg. USA). The output energy was measured with the joule metre, LED-200-C. Further details are given in appendix.

3.1.1 Optimization Of The Laser System

The e-beam laser output power depends on several parameters. These parameters were first optimized one by one by keeping all other parameters constant. Then the process was repeated to achieve the optimum performance of the laser. The laser system described in section (3.1) was optimized by varying the following parameters.

(1) Anode-Cathode Spacing Of E-beam Gun

The average energy of the secondary electrons was controlled by the magnitude of the imposed electric field of the sustainer bank and by the electron current coming out of the foil which in turn depends on the anode-cathode spacing of e-beam (see appendix). Both of them must be so adjusted that the electrons possess the optimal energy for efficiently pumping the gas molecules into the vibrational levels. So, for a fixed electric voltage of the sustainer bank; the anode cathode spacing which is the distance between the electron beam gun and the foil was optimized. For a fixed voltage of 42 KV at the sustainer bank, the anode cathode spacing of 100 mm was found to be optimum. It can be easily inferred from Fig.(3.8) as the laser output power was maximum in the case of Fig.(3.8f).

(2) Gas Ratio

Variations in the ratio of the gas mixtures changes the gain and the energy output of an e-beam laser due to various physical processes. The molecular kinetics (see appendix) such as vibrational energy transfer and decay processes are affected by the gas ratio. At a fixed value of average electron energy, changes in the gas ratio can produce large variations in the electron drift velocity / 63 %. Hence, optimum laser output was dependent also on gas ratio. The e-beam laser was optimized for various gas ratios, as shown in Fig.(3.9a). When He gas ratio was quite large, 80 or 90 %, then due to corona it was not possible to operate the e-beam laser with

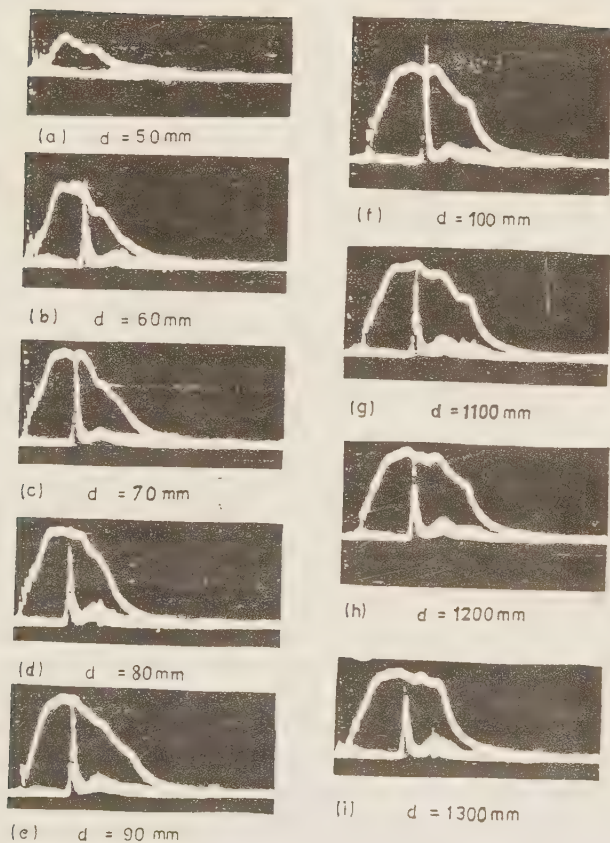


Fig.(3.8) The e-beam laser output power changes with the anode-cathode spacing, d . Bell shaped curve is the e-beam voltage and sharp pulse is the laser output. Maximum laser output power was available in case (f).

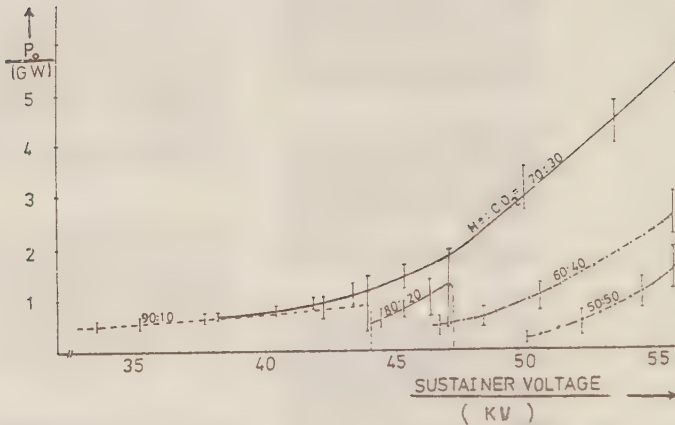


Fig.(3.9a) E-beam laser peak output power as a function of sustainer voltage for different gas ratios.

a sustainer voltage (E_s) more than 44 and 47 KV, respectively. On the other hand, small amplification was measured for a He gas ratio of less than 70%. It is obvious from the above figure that the optimum gas ratio is He:CO₂ = 7:3.

The output power increases with the sustainer voltage and a state could be reached when the laser is driven into the saturation state. It implies that a further increase in the input energy will not increase the gain and hence the output power of the laser will not increase anymore. Laser operation in the saturation state is undesirable as it not only reduces the efficiency of the system but also the laser pulse shape deteriorates. So, it was necessary to measure the output energy or power as a function of the sustainer voltage, as shown in Fig.(3.9a) and in Fig.(3.9b). Energy measurements were made for the purpose of comparison with the numerical modelling results as described in appendix. It shows that the maximum output energy for the optimum gas ratio was

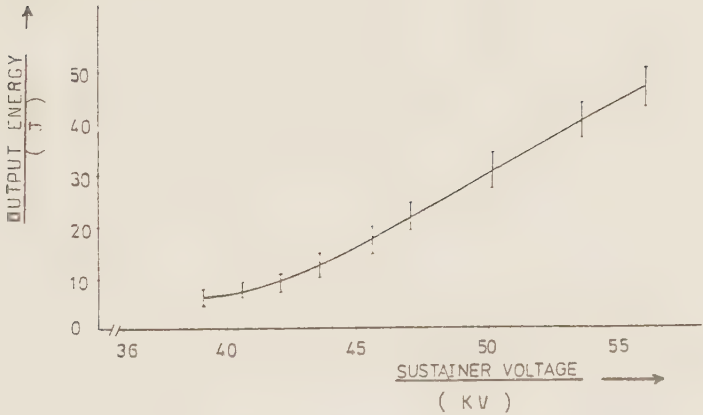


Fig.(3.9b) E-beam laser output energy as a function of sustainer voltage for the optimum gas ratio of He:CO₂ = 7:3.

about 45 J and the laser was not driven into the saturation state even up to a sustainer voltage of 56 KV, which was the maximum allowed voltage on the capacitors. The output energy of the e-beam laser was found not to be affected by IML.

Some of the above mentioned results have been confirmed by the computer modelling of our laser system as described in appendix.

3.1.2 Characteristics Of The High Power Laser System

Now, some important features of the high power CO₂ laser system will be described which are essential for a reliable and a clean experiment of LPI.

In the absence of IML, the e-beam laser output was found to contain P_{16} , P_{18} , P_{20} , and P_{22} lines, simultaneously, but predominated by P_{18} and P_{20} . However, it emits at a single line when good injection takes place.

A pulse of minimum of 6 KW from the master oscillator was found to be sufficient for IML in the slave oscillator. If all the losses due to various optics and SF_6 cell are taken into account, then the threshold for IML corresponds to an intensity of about 10^5 Wm^{-2} at the entrance of the e-beam laser.

The injected signal requires some time to be amplified in the slave oscillator. It must enter the cavity at the right moment, which is characterized by the term "locking window". It is the time interval during which if a signal is injected into the cavity, then the injection mode locking takes place, as illustrated in Fig.(3.10). Here, $\Delta t'$ is the time difference between the injected signal entering the e-beam laser and its emergence. It is difficult to recognize the individual pulses due to the large time scale and can be seen in Fig.(3.7c).

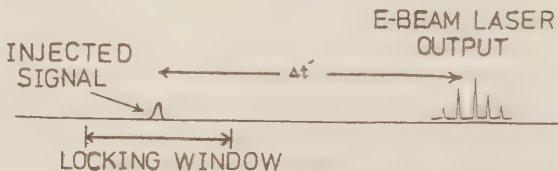


Fig.(3.10) (a) IML occurs when a laser pulse is injected into the e-beam resonator during the gain build-up, thus controlling the temporal behaviour of the laser emission.
 (b) Oscilloscope trace showing an injected signal and a train of pulses (shown inverted). $\Delta t'$ was about 450 ns.

The locking window was not found to be much critical rather than it was found to be as large as 600 ns. Moreover, no remarkable difference in the amplitude of the e-beam laser output pulses with the change in $\Delta t'$ within the locking window was measured. It therefore, resulted in a highly reliable laser output with IML as expected / 64 / and it was not difficult to achieve IML inspite of the inherent jittering of the master laser output.

Contrary to our observations, locking window was found to be short in other investigations /65-67 /. Belanger et al / 67 / observed a variation of about 25% in the output power due to the jittering in the injected signal of only 50 ns, while Alcock et al / 66 / reported a variation of 50% in the output within the locking window of 200 ns duration. Moreover, the mode-locked output power of the slave oscillator was found to be independent of the injected signal amplitude as long as the injected signal was above the threshold value. It resulted in the stability of the output power for a wide range of injected intensities.

The pulsewidth and shape of the e-beam output pulses follows very closely that of the injected signal, but the pulsewidth of the e-beam laser being somewhat shorter than the injected signal. No variation in the pulsewidth of the e-beam output could be found due to the change in $\Delta t'$ within the locking window range, contrary to the observations of Dyer et al / 68 /. Moreover, the effect of the structured pulse injection was also observed. E-beam output was found to follow the same structure as the injected pulse had. The e-beam laser output pulse has the same time difference between the spikes as that of the injected one inspite of different round trips within the resonator.

Hence, in this way, the master-slave CO_2 laser system has permitted the amplification of a few KW nanosecond pulse to power levels of a few gigawatt. So, a power density up to 10^{17} Wm^{-2} in the focal plane of 714 mm lens can be achieved. In one shot of the laser output, there was not only a single

pulse, but more than 10, with an interseparation time of 16.7 ns. The laser output frequency distribution had a single axial mode due to IML.

A laser with above mentioned characteristics was found to be appropriate for the study of parametric instabilities.

3.2 The Z-Pinch

As already stated, a laser plasma interaction experiment was intended to be performed under well defined conditions. If the same laser is employed for the production and also for the diagnosis of the plasma, then the motion and inhomogeneities in the plasma develop which can yield misleading results. So, it is essential for the investigation to have a preformed and well diagnosed plasma.

Moreover, there are many advantages to study a gas target over solid targets. They are:

- a homogeneous and stationary plasma of a variable density can be produced. It can span the whole range of plasma densities from a critical to well below critical density,
- absence of the critical surface formation (for $n < n_{cr}$) eliminates the classical reflection from a moving critical surface inside the plasma column,
- long interaction length makes possible to achieve the theoretical limits like 100% back scattering, and
- nondestructive target which can be synchronized with the laser.

In the experiment reported here, the plasma was produced in a dynamical Z-pinch device. The discharge chamber was constructed in a cylindrically symmetric manner using a quartz vessel. It has an inner diameter of 180 mm, with stainless-steel electrodes about 500 mm apart. Midway between the electrodes and diagonally opposing one another are four ports, for diagnostic purpose.

The bank powering the discharge device consists of nine 4.5 μ F capacitors charged to 32 KV and it delivers a total discharge energy of 20 KJ. Directly at the top of each capacitor was a spark gap switch, as shown in Fig.(3.11). The total inductance

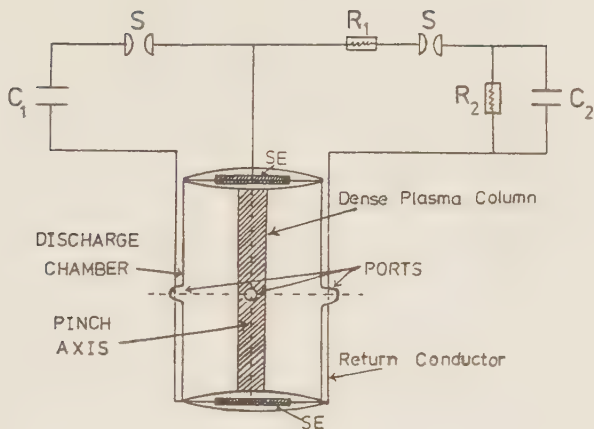


Fig.(3.11) A schematic diagram of the Z-pinch device, where C_1/S = one capacitor of $4.5 \mu\text{F}$ / spark gap arrangement out of 9 such circuits switched in parallel,
 $C_2 = 2.2 \mu\text{F}$ capacitor for preionizing discharge,
 $R_1/R_2 = 0.2\Omega / 25\Omega$, and
 SE = stainless steel electrode.

of all passive elements in the discharge circuit was , thus, kept at 30 nH.

This dynamical Z-pinch device was set up and diagnosed by Finken et al /69-71 /. It was found that the gas filling pressure can be varied in the range of 33 - 665 Pa resulting in a plasma density of $0.8 \times 10^{24} - 5 \times 10^{24} \text{ m}^{-3}$, i.e., from 0.05 to $5 n_{\text{cr}}$. Mostly, the experiment was performed at a plasma density of $0.16 n_{\text{cr}}$ since an underdense plasma was desired and the plasma parameters were well known for these conditions. At this density and at the time of maximum compression, the temperature $K_B T_e \approx K_B T_i = 10 \text{ eV} \pm 10\%$ and density were reasonably

homogeneous, the H_2 plasma was fully ionized with $Z_{eff} \approx 1$, and the dense plasma column had a diameter of 18 mm. The radial density distribution as shown in Fig.(3.12) persists for about 80 ns about the time of maximum compression with dynamical effects preceeding and following. Further information about the Z-pinch is given in Fig.(3.12) and in the specifications in appendix.

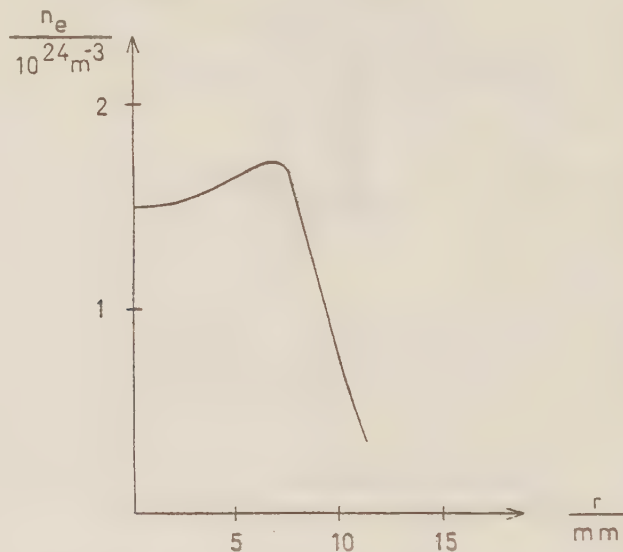


Fig.(3.12) Radial density distribution of the plasma column in Z-pinch / 70 /.

3.3 The Scattering Experiment

In the scattering experiment, an homogeneous underdense H_2 plasma (section 3.2) was irradiated with a high power CO_2

laser system (section 3.1). Mainly, the

- forward scattering, and
- backward scattering of the electromagnetic radiation due to SBS was investigated in detail.

3.3.1 The Forward Scattering Experiment

The experimental set up for the forward scattering experiment is shown in Fig.(3.13). The scattered laser light in the forward direction was investigated both temporally and spectrally. The spectra of forward scattered or transmitted laser light were searched near ω_0 to look for, the forward scattering, normal transmission, and multiple scattering.

Short CO_2 laser pulses were focussed onto the preformed Z-pinch plasma by a 1.5 m focal length salt lens. Only NaCl optic was used as it has a low absorption coefficient, low refractive index resulting in a small reflection of about 4% from each of its surfaces and it is not expensive.

High damage threshold mirrors M_1 and M_3 were employed for directing the laser beam over long distances to avoid late pulses from e-beam laser to reach the Z-pinch plasma at the time of its maximum compression. Two wedge shaped beam splitters were employed for monitoring a fraction of the incident and the scattered radiations without any interference between them, without reducing the laser intensity to a significant amount and to reduce the optical components to a minimum. As shown in Fig.(3.13), the main laser beam directed by M_3 passes through the beam splitter BS_1 , directed by M_1 , and finally focussed by a plane convex lens f_1 of 0.7 m or 1.5 m focal length on the axis of the dense plasma column in the Z-pinch. The first reflection from BS_1 was used for monitoring the laser radiation. This reflected light passed through another beam splitter BS_2 and was further divided. The front surface reflection of BS_2 after two reflections from M_9 and M_8 was focussed by a 0.7 m focal length lens onto the detector D_1 . The cross-section of the detector was large enough to record the full intensity of the laser. It not only gives the information

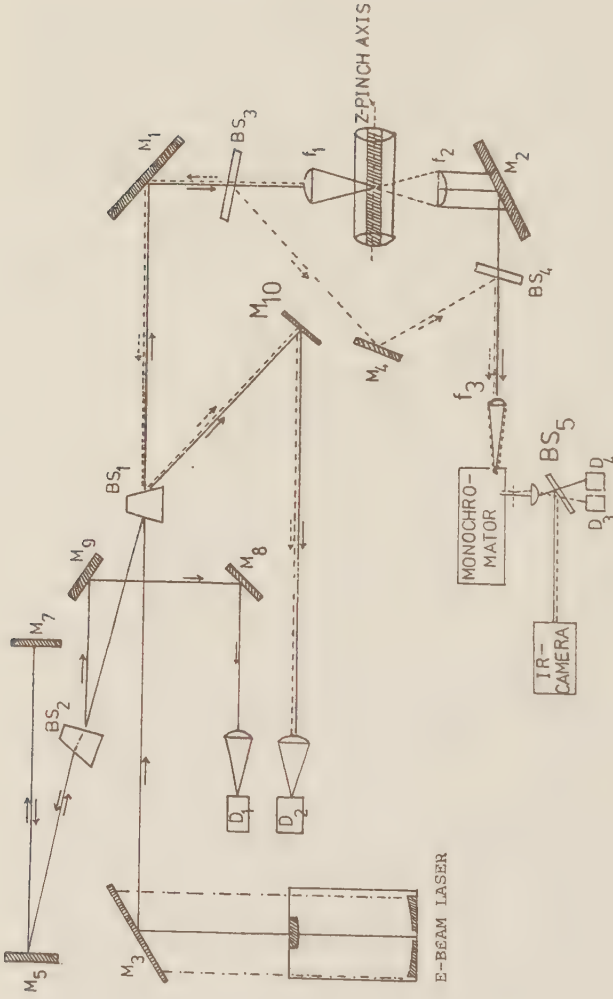


Fig.(3.13) The schematic experimental arrangement for the investigation of temporal and spectral aspects of SBS both in the backward and in the forward direction.

about the power of each individual pulse but also about their shapes and eventual finite background radiation in between single pulses. Another fraction of the incident laser radiation passes through BS_2 , then back and forth between M_5 and M_7 mirrors in order to achieve an optical delay necessary for recording the incident and the scattered laser pulses at one and same oscilloscope, after that it passes through BS_2 and BS_1 again, reflected by M_{10} and is finally focussed by a 0.7 m focal length lens at a photon drag detector D_2 . At the same time, a part of the radiation backscattered into the solid angle of f_1 from the plasma was reflected by the front surface of the beam splitter BS_1 . It was directed by M_{10} and was focussed at the same detector D_2 . D_2 has a sensitive area of 200 mm^2 which is sufficient enough to record all the radiation. In this way, scattered and incident radiation are detected with the same sensitivity and rise time. So, it reduces the error in the determination of the amount of backscattering to a minimum. The radiations were not focussed exactly at the detectors D_1 and D_2 to avoid any plasma formation and to prevent the detectors from being damaged. In addition, special care was taken that the radiation did not fall out of the sensitive area of the detectors.

Another part of the backscattered radiation was provided by the front surface reflection of BS_3 . It was directed by the mirror M_4 and beam splitter BS_4 to the IRSA system for spectral analysis of the backscattered radiation. A 1.5 m focal length lens finally focusses the light at the entrance slit of the monochromator. Lens f_2 was employed to make the emergent light parallel (coming out of the Z-pinch). It was also focussed by the same lens at the monochromator. The pinch exit port with 50 mm diameter was sufficiently wide for all near 180° scattered light to pass. The emerging radiation might have been forward scattered, multiply scattered, or been transmitted one. So, a monochromator was employed to separate them. Two separate detectors D_3 and D_4 were installed to provide time resolution and to record the intensity of two selected wavelength regions. In this way, a time delay between

scattered and incident laser radiation can also be determined besides the frequency shift. The transmitted light has the frequency of the incident laser but the scattered light was expected to have some frequency shift near ω_0 . To increase the spatial separation of the two detectors an optical system with 1:3 magnification was employed behind the monochromator. To provide the readings of time integrated spectra, a beam splitter BS₅ was employed to direct a fraction of the transmitted and scattered laser light to the IR camera target.

Mostly, optical components of 110 mm were employed as the full cross-section of the laser beam was 100 mm in diameter. Reflections from the surfaces of the wedge shaped beam splitters reflect a fraction of the light. So, the incident and scattered radiation can be monitored by various detectors. Three photon drag detectors, D₁, D₂, D₃ and a pyroelectric detector D₄ were employed to record the intensity of the incident and scattered radiation and for the high temporal resolution, which was achieved by coupling D₄ to a fast Tektronix 7104 oscilloscope with 7A29 plug in unit. The combined rise time of the system was about 350 ps.

An infra red system analyzer (IRSA) was employed for the spectral resolution of the scattered radiation. The spectra recorded by IRSA is a time integrated one. It consists of a 1.5 m 1:8 Czerny Turner monochromator to which an Hamamatsu IR-TV system was coupled for electronic and digitized read out. The overall spectral resolution of the IRSA was better than 0.6 GHz at 10.6 μm .

Calibrated absorbers were employed for attenuating the beam. The time correlation between the laser and plasma at the time of interaction was determined by recording the photomultiplier signal of the Bremsstrahlung radiation from the plasma and the laser monitor signal, simultaneously. The peak of the Bremsstrahlung signal corresponds to the time of the maximum compression of the plasma column in Z-pinch, as shown in Fig.(3.14).

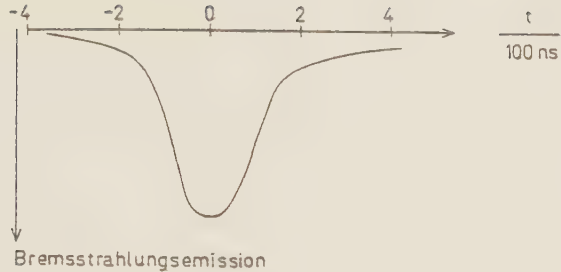


Fig.(3.14) Bremsstrahlung emission signal of the plasma as a function of time.

The IRSA system was wavelength calibrated with a CW CO_2 tunable laser, while the detectors D_1 , D_2 etc. were calibrated relative to each other by directly irradiating them with the laser beam. Also is shown the geometry of the irradiated Z-pinch plasma column in Fig.(3.13)

3.3.2 The Backward Scattering Experiment

In another experiment also the backscattered radiation was investigated. An attempt was made to resolve temporally the backscattering spectra which to date, had been measured time integrated only. The intention was to gain in this way some information about the saturation mechanisms.

For this purpose, it was essential to modify the experimental set up described in the previous section (3.3.1). In

the modified experimental arrangement, as shown in Fig.(3.15), again special care was taken to reduce the number of optical components to a minimum. All the optical components (f_2 , M_2 , BS_3 , and BS_4) installed for the detection of forward scattering, as shown in Fig.(3.13), were removed. Moreover, the detectors D_3 and D_4 were replaced by an array detector for the temporal resolution of the backscattered light spectrum.

A portion of the radiation backscattered into the solid angle of the focussing lens f_1 was reflected by BS_1 and was focussed by a 1.5 m focal length lens f_3 onto the entrance slit of the monochromator. A fraction of the incident radiation was provided by BS_2 , directed by mirrors M_5 , and was finally focussed by the same lens f_3 onto the entrance slit. In this way, both the incident and backscattered signals were spectrally resolved by the monochromator, simultaneously.

For a better resolution, the spectrally resolved light emerging out of the monochromator's exit plane was reconstructed by a 10cm focal length lens with 1:3 magnification at the camera target and at the array detector, as shown in Fig.(3.13). The camera target has 16 mm diameter and has a nearly constant response in the $8\mu\text{m}$ to $14\mu\text{m}$ range. The IRSA system has again a spectral resolution of about 0.6 GHz in this experimental set up as before. The camera target position was checked to be aligned to best image by replacing the diffraction grating in the monochromator with an aluminium mirror and by recording the image of the entrance slit at the camera target. The minimum slit width was taken to be the best image position.

The camera diameter was sufficiently wide to cover both the wavelength position of the incident and of the backscattered radiation and in this way permitted to determine directly the wavelength shift of SBS. On the other hand, array detector was so installed that only the backscattered light was resolved on its six channels. Thus, the IRSA system records a time integrated spectrum of the backscattered radiation of all pulses in one shot, while the array detector resolves also in time the spectrum of each individual backscattering pulse.

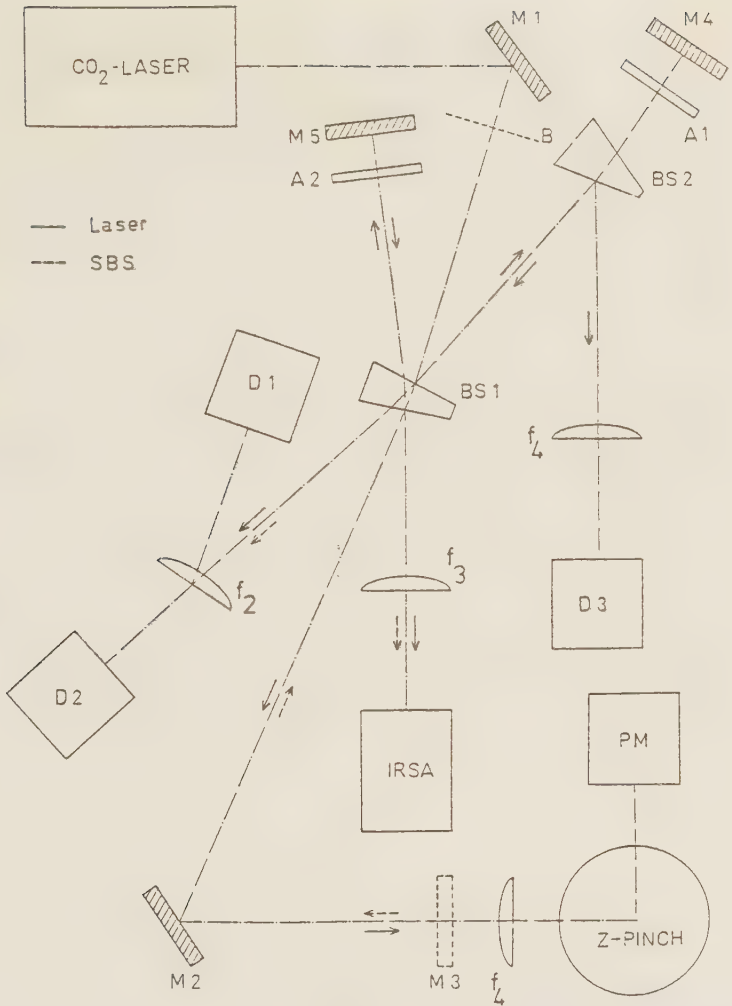


Fig.(3.15) Experimental arrangement for the temporal resolution of the backscattered light spectra.

The array detector consists of six UF_1 detectors in a row, each 0.8 mm wide with 1 mm distance between the centres of the detectors. The UF_1 detector of Pulse Systems (USA) is the fastest available pyroelectric detector in the market. It has a rise time of 50 ps with a sensitive area of 1 mm^2 . In this experimental set-up, each detector covers a spectral range of 1.8 GHz. All the array detector signals were recorded with the fastest Tektronix oscilloscope 7104. These signals were so adjusted that they fall in between the two backscattered pulses on the oscilloscope screen. Therefore, cables with different lengths were used for each detector and only five detection channels were employed in order to avoid partial overlapping. The time difference between the signals and the relative sensitivities of the detectors were calibrated by directly irradiating them with the laser radiation.

As before, the detector D_3 has provided information on the time dependence of the incident laser power, while D_2 has yielded information about the amount of the backscattering which can be easily calculated from Fig.(3.16). For the high temporal resolution of the backscattered pulses, a UF_1 detector indicated here with D_1 was employed. The geometry of the laser plasma interaction experiment was the same as shown in Fig.(3.13).

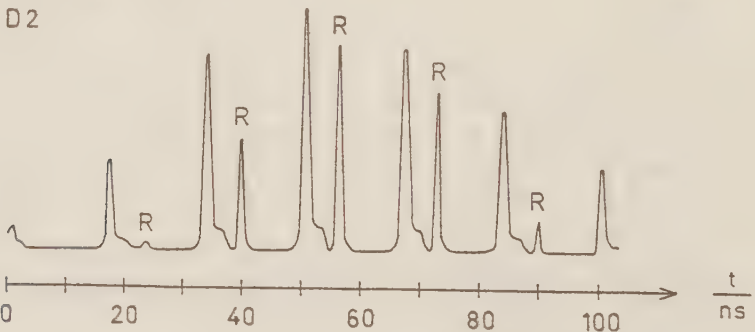


Fig.(3.16) Incident and the backscattered pulses (indicated by R) as recorded by detector D_2 .

The experiment was performed at $0.16 n_{cr}$ as at this density the plasma parameters were known and SBS had been investigated at this density. /16,31 /.

Chapter 4

Results and Discussion

In this chapter the results of the scattering experiments will be described. These results will be analyzed in the light of the theories presented in the second chapter. Accordingly chapter 4 is divided in three main sections. In the first section, processes other than SRS are dealt with while in the next two sections the results of the SRS scattering experiments in forward and in backward direction will be described and discussed.

4.1 Parametric Decay Processes

In the experiment described, provision was made to vary the target plasma density from 0.08 to $5 n_{cr}$ (section 3.2). So, an attempt was made to observe SRS, SCS, and SRS in different density regions.

4.1.1 Stimulated Raman Scattering (SRS) and Stimulated Compton Scattering (SCS)

The SRS experiment was performed at the low plasma density of $8 \times 10^{23} \text{ m}^{-3}$ which was much below the highest density $2.5 \times 10^{24} \text{ m}^{-3}$ for SRS to occur. This was necessary in order to shift the SRS wavelength to the observable wavelength range below. Radiation above $15 \mu\text{m}$ could not be detected due to the limited transmission bandwidth of NaCl optics. However, this does not represent an undue restriction of the experiment as it has been recently observed / 72 / that SRS occurs at electron densities much lower than the above limitation.

The experimental arrangement was that described in section (3.3.2) with the 1.5 m monochromator replaced by a 0.25 m Jarrel Ash monochromator. The IR-camera system was an order of magnitude more sensitive than the detector D_2 . The incident laser radiation was attenuated but not the backscattered one. In order to detect also very weak SRS radiations, the sensitivity of the detection system was made higher by a factor 10^3 as compared to the SRS. Both the

laser intensity and the plasma density were varied. In spite of all efforts, SRS higher than 0.1% of SRS in the range of 10-15 μm could not be detected.

The reasons could be: the direct dependence of the SRS frequency causes the wavelength position to change from shot to shot in a relatively wide range which could not be covered by the detectors. This density change was caused by irreproducible functioning of the Z-pinch device at the required low densities which could not be controlled. It was, anyhow, the lower limit of the plasma density which could be produced in this Z-pinch.

Although the intensity of the laser was much above the SRS threshold of $3 \times 10^{12} \text{ Wm}^{-2}$ (section 2.1.2), the e-folding time for growth was about 500 ps. For a laser pulse of 2 ns (FWHM), it might not have been sufficient for SRS to grow to detectable levels.

The threshold of SCS (section 2.1.1) also was only 10^{13} Wm^{-2} for the plasma investigated. But no SCS could be measured, neither in the forward (section 4.2.3) nor in the backward (section 4.2.2) direction. We also interpret this result as a consequence of the short laser pulses in combination with the experimentally established fact that the temperature ratio T_e/T_i increases rapidly above unity at the beginning of the pulse. In this way, the growth rates are decreased and during the short laser pulse there is not sufficient time left for SCS to attain scattering levels which were detectable.

The absence of SRS and SCS have paved the way for high amount of SRS.

4.2 Stimulated Brillouin Scattering (SBS)

An homogeneous fully ionized hydrogen plasma of 0.16 n_{cr} density was irradiated with a laser of maximum intensity of 10^{17} Wm^{-2} . SRS was investigated in the experimental set up shown in Fig.(3.13) and in Fig.(3.15). The scattered radiation both in the forward and in the backward direction was observed. The scattered radiation was both temporally and spectrally resolved with

the fast detection system and with the IRSA system, respectively, as described in section (3.3). First of all back-scattered radiation will be investigated which will help to understand the nature of the forward-scattered radiation.

4.2.1 Backscattered Radiation

To gain some insight into the nature of the scattered radiation, it is necessary first to find out the threshold and the maximum amount of backscattering.

As described in section (3.3) and shown in Fig.(3.7c), detector D_1 yields the temporal information about the incident laser pulses within a shot. In case of good IML, all the laser pulses have the same shape and pulsewidth within a shot, but only the power varies from pulse to pulse. So, up to ten pulses of a single shot span a wide range from threshold to a significant amount of the backscattering.

For the measurements of the amount of backscattering, detector D_2 was employed with a typical laser shot shown in Fig.(3.16). SBS due to the first laser pulse was below than our detectable level but the SBS due to the second laser pulse could be observed which implies that the threshold of SBS lies between the intensity of the two laser pulses. The error in the determination of the backscattering was small as both signals were recorded by the same detector. So, the relative error in the determination of the backscattering was estimated to be about 10%.

The relative backscattering values P_B/P_0 of all pulses in several shots occurring within ± 40 ns about the time of maximum compression of the plasma in Z-pinch have been plotted as a function of the peak intensity of the corresponding laser pulse, Fig.(4.1). A time range of ± 40 ns was selected for the laser interaction with the plasma as it was known {section (3.2)} that the plasma remains stationary in this time range. The purpose was to avoid any dynamical effects which will complicate the situation.

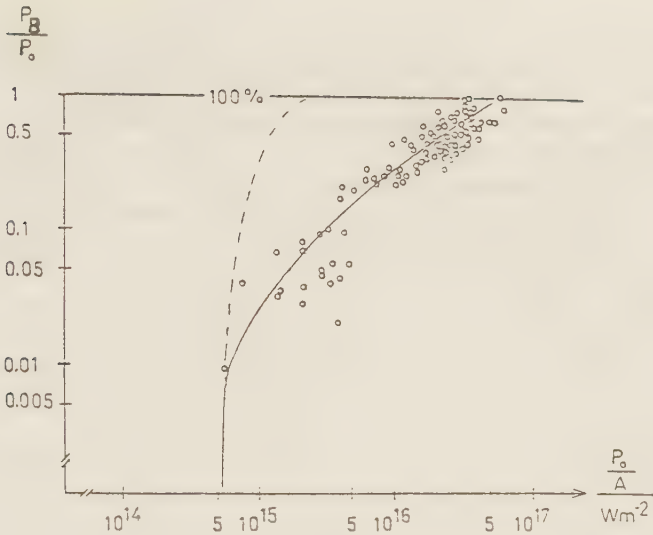


Fig.(4.1) Relative back scattering power as a function of laser intensity in several shots covering the whole range from threshold to 100% for a plasma density of $0.16 n_{cr}$. Laser focal spot area, A , was 0.07 mm^2 . P_B and P_0 represents the back scattered and incident power. Solid line is the best curve fitting, while dotted line is the theoretical limit.

The measurements span the whole range from threshold to 100% backscattering. By extrapolating the backscattering curve, the threshold laser intensity in the plasma at $0.16 n_{cr}$ was found to be $5 \times 10^{14} \text{ Wm}^{-2}$ for $A = 0.07 \text{ mm}^2$, where A is the area of the laser focal spot in the plasma. The effect of the density variation on the threshold was also measured. It was observed that the threshold decreases with an increase in the plasma density. The threshold of SBS was found to be in good agreement with the theoretical predictions for the homogeneous plasma and has been discussed in detail by Handke /31/. It confirms the nature of the backscattered radiation as SBS.

Plotting the backscattering of only the various pulses in a single shot yielded results as shown in Fig.(4.2). These data points show considerably less scatter and a straight line can be drawn through the points. This means that the larger scatter observed in Fig.(4.1) must be traced to shot by shot differences in the incident laser pulses or the plasma conditions.

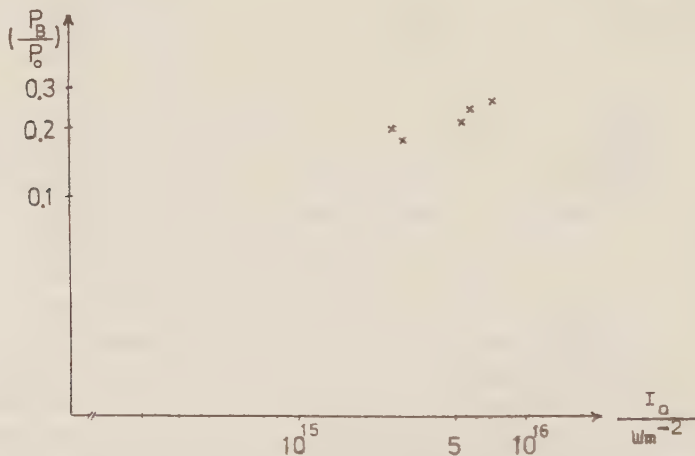


Fig.(4.2) Relative backscattering power as a function of laser intensity within a shot (#39301).

The most important feature of this experiment was the observation of about 100% backscattering. It does not imply that no saturation mechanism was present. From the shape of the curve, it is obvious that some mechanism has limited the rapid growth of SBS beginning at about 10%. It was one of the main objectives of this investigation to find a clue as to the nature of this anomalous damping mechanism. The idea was to gain insight into this unknown mechanism by studying the back scattered light spectrum both time integrated and time resolved.

4.2.2 Time Integrated Backscattered Light Spectrum

Although the frequency shift ($\Delta\nu$) of the back scattered light was only in the vicinity of 20 GHz it could be resolved accurately with the IRSA system. When the backscattered light was spectrally resolved, it was found to have characteristics different from that of the incident laser light. Two typical spectra of the incident and backscattered light near and far from threshold, i.e. near 100% backscattering are shown in Fig.(4.3). Incident laser light was much more attenuated than the backscattered one. The study of the backscattered light spectra shows the following important features:

- backscattered light spectra near threshold and during the stationary plasma phase has a frequency shift to the red of about 17 GHz relative to the incident laser light,
- scattered light had a spectrum only slightly broader than the incident one,
- shift and width of the scattered light spectrum increases very slowly with the amount of backscattering which in turn depends on the laser intensity,
- frequency shift increases from 17 GHz near threshold as shown in Fig.(4.3a) to 23 GHz, Fig.(4.3b), near 75% backscattering,
- if the laser has not interacted with the plasma at the time of maximum compression, but rather quite early or late during the dynamic plasma phase, then correspondingly less or more frequency shift was observed due to Doppler effect. But the net frequency shift was always towards red as shown in Fig.(4.4),
- no blue shifted component was observed,
- both frequency shift and width of the scattered light spectrum increases for the same amount of backscattering, if the laser radiation does not drop to zero in between subsequent pulses of one shot,
- backscattering spectra had reproducible features, and
- no rescattered light at 3ω was observed

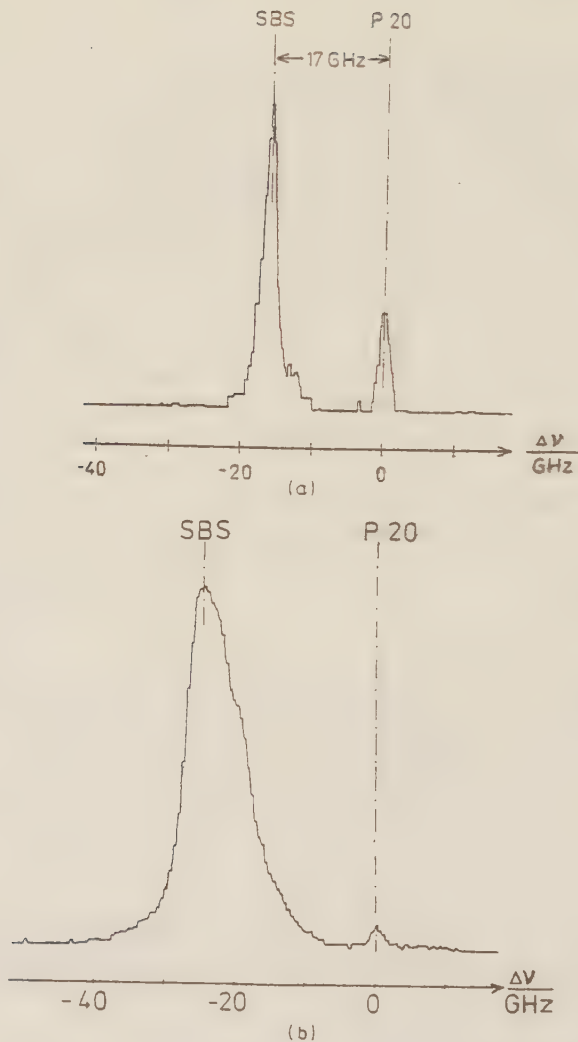


Fig.(4.3) Time integrated spectrum of the backscattered and the incident laser light as recorded by the IRSA system for the case of a
(a) low backscattering (about 10%), and
(b) large backscattering (about 75%).

The above mentioned features help to identify the nature of the backscattered light as there are many possibilities for the laser light to be reflected.

The observed frequency shifted component could not have been due to the geometric reflection of the laser radiation from a critical surface, as the plasma density was much below that of critical density. Moreover, a red shift in the backscattered light spectrum was always observed even if the laser was fired before or after the maximum compression of the plasma, i.e. when the plasma was not stationary. If the frequency shift would be due to the bulk motion of the plasma then both red or blue components should have been observed. Further confirmation comes from Fig.(4.4) in which the observed frequency shift in the backscattered radiation versus the time difference between the incident laser and the plasma maximum compression for a constant laser intensity

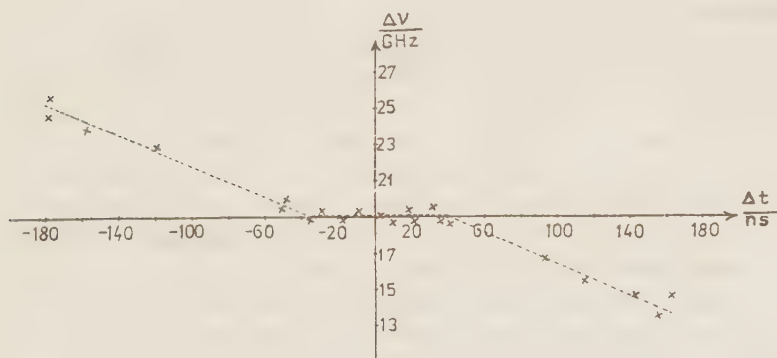


Fig.(4.4) Frequency shift (ΔV) in the backscattered radiation relative to the incident laser radiation as a function of the time coincidence (Δt) between the laser maximum and the maximum compression of the plasma for a fixed laser intensity of 10^{16} W m^{-2} .

is plotted. It shows that the plasma was stationary about the time of its peak compression for $\Delta t = \pm 40$ ns as in this time range a constant frequency shift was always observed for the same laser intensity. It is in good agreement with the experimental observations of Finken et al /69-71/ on this pinch. From their measurements, it is known that the plasma was moving with a speed of 4×10^4 ms⁻¹ at a time of $\Delta t = 100$ ns or $\Delta t = -100$ ns corresponding to a Doppler frequency shift of 3.8 GHz, which is in good agreement with Fig.(4.4). If the laser has interacted with the plasma when it was not in the phase of maximum compression then the dynamical effects can be observed in the backscattered radiations. The frequency shift increases or decreases for $\Delta t < -40$ ns or $\Delta t > 40$ ns, respectively.

The above considerations show that the backscattering must have been due to some parametric processes in the plasma. The possibility of the modulational instability can be ruled out on the ground that only one side band, the low frequency component, was observed. Moreover, the low frequency shift rules out the possibility of parametric decay processes other than SRS and SCS as the non-linear mechanism responsible for the laser light scattering. SRS can be ruled out on the ground that the observed frequency shift was quite small.

If the backscattering was due to SCS, then the observed frequency shift would be $\Delta \omega = 2 k_o v_j$, where v_j is the ion or electron thermal velocity and k_o is the wave number of the incident laser radiation. In case of SRS, 17 GHz frequency shift near threshold corresponds to the weighted sum of electron and ion temperature of 100 eV. But the electron and ion temperature could be determined / 31 / with the help of a heating model to be $K_B T_e = 52$ eV, and $K_B T_i = 16$ eV. It implies a frequency shift of about 9 GHz for ion SCS and about 700 GHz for electron SCS which is obviously at variance with the observations. So, SRS is the only mechanism left.

All the above analysis shows that only SRS was responsible for the laser light scattering in the backward direction. However, it will be further confirmed by observing the nature of the scattered radiations in the forward direction.

4.2.3 Forward Scattering

Here, the results of the forward scattering experiment will be discussed. The experimental set up has already been described in section (3.3.1). Most of the forward scattering experiment was performed at a plasma density of $0.08 n_{cr}$. Such a low density was selected so that the laser light could be transmitted easily through the plasma. Altogether four detectors were employed, D_1 yields the information about the incident laser pulses such as pulse power, its duration, shape and background, if any. D_2 records the monitor and the back scattered signal, simultaneously. D_3 measures the transmitted laser light and D_4 has recorded the frequency shifted light in the forward direction. In addition to these, Bremsstrahlung signal from the plasma, and the monitor signal from D_1 yield the information about the time correlation between the laser and the plasma. Moreover, a time integrated spectrum was also recorded with the help of the IRSA system. Some red frequency shifted light was observed in the forward direction along with the transmitted laser light. So, an attempt was made to explore the nature of this frequency shifted light in the forward direction by studying the following aspects:

- Threshold

First of all it was necessary to find out the threshold value of the forward scattered (FS) radiation which will help in its identification. For the purpose of comparison, the threshold for back scattering will be also discussed here. For each laser pulse (above the threshold) the back scattered signal as viewed by D_2 was seen on the oscilloscope screen about 6 ns after the corresponding laser pulse. Much of the FS radiation was absorbed in the plasma and hence the FS signals were quite weak at the detectors. But the camera was an order of magnitude more sensitive than the detectors and was appropri-

ate for the forward scattering threshold measurements. It was found that the camera had recorded some frequency shifted light in the forward direction only when back scattering was present. By taking into account all the losses, the threshold for FS and SBS can be determined.

Threshold for FS was found to be $6 \times 10^{15} \text{ Wm}^{-2}$ for $n_e = 0.8 \times 10^{24} \text{ m}^{-3}$, and $A = 0.07 \text{ mm}^2$, where A was the area of the laser focus in the plasma. For the same experimental conditions, the threshold for the back scattering was measured as $4 \times 10^{15} \text{ Wm}^{-2} / 16$. So, both the forward scattering and back scattering had the same threshold within the experimental (error) limits.

- Frequency Shift

There were three possibilities of the scattering processes in the forward direction near ω_0 . The first one was SBS in the forward direction almost with the same frequency as in the back-ward direction. Another possibility was the rescattering process. If the length of the interaction volume is sufficiently large then it is possible for the back scattered light to be scattered again. The rescattered light in the forward direction must have twice the frequency shift as compared to SBS back scattering and the third possibility is SCS. For the identification of the scattered light in the forward direction, it is necessary to first summarize the experimental facts related to it.

A definite prove for the identification of a parametric process is the frequency shift of the scattered radiation relative to the incident radiation. Spectra of FS radiation were recorded by the IRSA system described in section (3.3.1). A typical spectrum of the incident and FS radiations is shown in Fig.(4.5). The laser was fired into the plasma at the time of its maximum compression. So, the plasma was stationary and

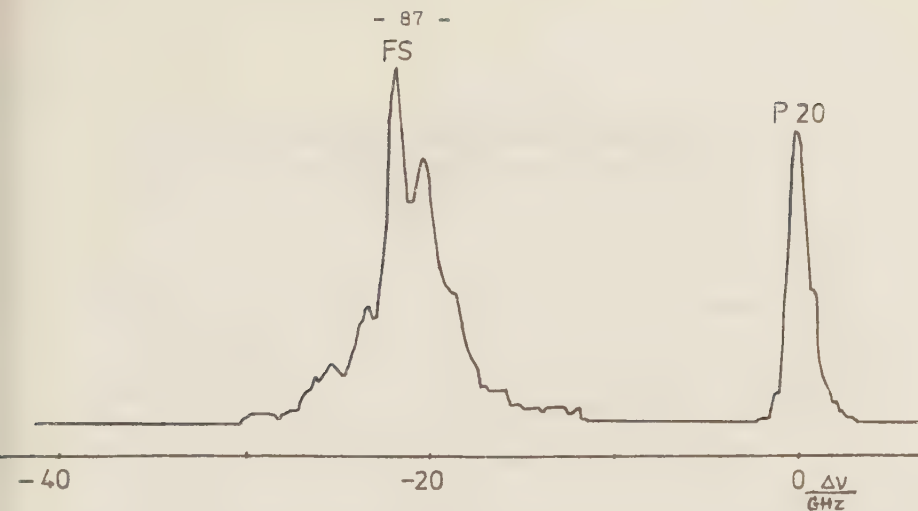


Fig.(4.5) Time integrated spectrum of the incident laser light (P 20) and the scattered light (FS) in the forward direction. FS spectra is frequency shifted to the red by 21 GHz and is broadened. The maximum amount of backscattering in this shot was about 50%,
 # 33936.

there was no chance for the additional frequency shift to be due to a Doppler effect. In this shot, the back scattering was 50% corresponding to laser intensity of $3 \times 10^{16} \text{ Wm}^{-2}$. The FS light has a spectrum much broader than the incident laser light with a red shift of about 21 GHz. This frequency shift and the broadening features of the FS spectrum were identical with that of the back scattered light spectrum (section 4.2.2) under the same experimental conditions. No difference could be recognized. Because of the 21 GHz-red shift the possibility of SCS with its frequency shift of 9 GHz must be ruled out.

An attempt was made to check the rescattered light in the forward direction at twice the frequency shift. No signal up to a detection level of 10^{-3} SBS intensity could be

observed. On this basis, the possibility of rescattering in the forward direction to a significant amount could be ruled out.

- Temporal Coincidence

No remarkable difference between the characteristics of forward scattered and back scattered light due to SBS could be established. To resolve the paradox, the time difference between the transmitted laser light through the plasma and the FS signal was exactly determined as shown in Fig.(4.6). It amounted to 350 ns. As was finally found, this time delay corresponds very well to the SBS-signal, travelling back to the e-beam resonator, being reflected there and again irradiating the plasma at an instant when it has become better transparent, thus, being now recorded in the forward direction.



Fig.(4.6) Oscilloscopic trace of the incident laser light (downward) and the FS light (upward) . Horizontal scale was 100 ns/ division.

4.3 Time Resolved Backscattered Light Spectrum (TRBS)

The above mentioned experimental results confirm that only SRS is occurring in the backward direction and is responsible for the large amount of backscattering and there is no forward scattering. But it does not help to understand a large deviation of the backscattering curve, which becomes apparent near 10% SRS, from the theoretical one as shown in Fig.(4.1). Besides this, an increase in the laser intensity above threshold was found resulting in a small increase systematically in the frequency shift towards red and hence the plasma temperature increases correspondingly. A highly satisfactory explanation on the basis of a simple heating model has been worked out by Handke / 31 /. His model predicts $K_B T_e = 52 \text{ eV}$, $K_B T_i = 16 \text{ eV}$ near threshold and $K_B T_e = 108 \text{ eV}$, $K_B T_i = 19 \text{ eV}$ near 100% backscattering. He gave a relationship for the way the SRS-red shift should vary with incident laser intensity. According to this model, the redshift increases slowly with the laser intensity as $I_0^{1/7.5} \propto \Delta \nu$. In addition to the other predictions of the model, he had also checked in the case of a single high power pulse that also the redshift prediction was right. However, with the increase in the amount of backscattering, the spectra get broadened and it impedes to associate the laser pulses with the peaks of the backscattered light spectra within a shot. It makes difficult to check the $I_0^{1/7.5}$ law. On the other hand, the Forslund model /8/ in the strong coupling region predicts $I_0^{1/3} \propto \Delta \nu$.

Since the theory of Forslund predicts a different dependence in case of strong coupling and occasional time-integrated SRS-spectra indicated deviations from the $I_0^{1/7.5}$ law, this aspect was investigated in detail with a large number of SRS-spectra in a wide intensity range. The most informative data that could be taken for a deep insight into the subject in this experiment would be single pulse and time resolved back scattered light spectrum. The experimental procedure used for recording TRBS has already been described in section (3.3).

Time integrated spectra were recorded for all the pulses in a shot with an IRSA system, which has now 3.2 magnification optics behind the monochromator. The 0-order reflection from the grating was recorded by the IRSA system to check the time-integrated spectra at greater spectral resolution. The apparatus profile of the IRSA system was found to be limited by the camera target to 0.6 GHz.

The time resolved spectra of each individual pulse within a shot were recorded with the help of an array detector, AD. A typical shot with the backscattered pulses along with the corresponding spectrally resolved signals from the different channels of the array detector is shown in Fig.(4.7). All channels had different cable lengths to avoid any overlapping of the signals. These signals correspond to different frequencies present in the backscattered pulse. The amplitude

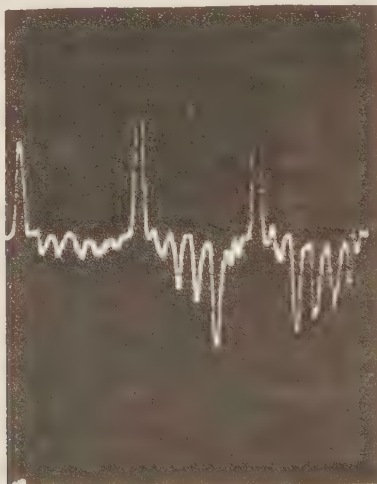


Fig.(4.7) Backscattered laser pulses (upwards) as viewed by the detector D_2 and the spectrally resolved signals (downwards) from the different channels of the array detector.

of the signal represents the intensity of the SBS light in a particular frequency range of 1.84 GHz which corresponds to the spectral bandwidth of a channel in the array detector of this experimental set up. On the other hand, detectors D_1 and D_2 had yielded again the information about the incident laser pulses and the corresponding SBS reflectivity, as in the previous experiments. The error in the frequency and time determination was ± 0.9 GHz and ± 100 ps.

A study of large number of time resolved spectra has yielded the following important results:

- In principle, if the time resolved spectra corresponding to all the laser pulses in a shot are superimposed, then it must reproduce the time integrated spectrum. A typical shot is shown in Fig.(4.8). In this shot three main laser pulses were temporally resolved. It was found that, if the backscattering is less than 10%, then a single pulse spectrum has a peak near 17 GHz and has a small bandwidth of less than 4 GHz. It is in good agreement with the previous observations (section 4.2.2), where the backscattered light spectra near threshold were observed to be red shifted by 17 GHz. Moreover, it can be easily seen in Fig.(4.8) that the time resolved spectrum reproduces all the essential features of the time integrated spectrum.
- In the absence of any background heating, if the amount of backscattering was less than 10%, then single pulse spectra were not splitted. Otherwise, they show two peaks. The first observed backscattered pulse in a shot had a reflectivity less than 10% and its spectrum (--- curve in Fig.4.8) is not splitted, while the other two had reflectivities of about 15% each and their spectra are splitted.
- For backscattering greater than 10%, saturation effects are obvious not only in the spectra but also in the different wavelength channel signals as shown in Fig.(4.9). Some of the signals corresponding to the splitted spectra were even modulated in time.

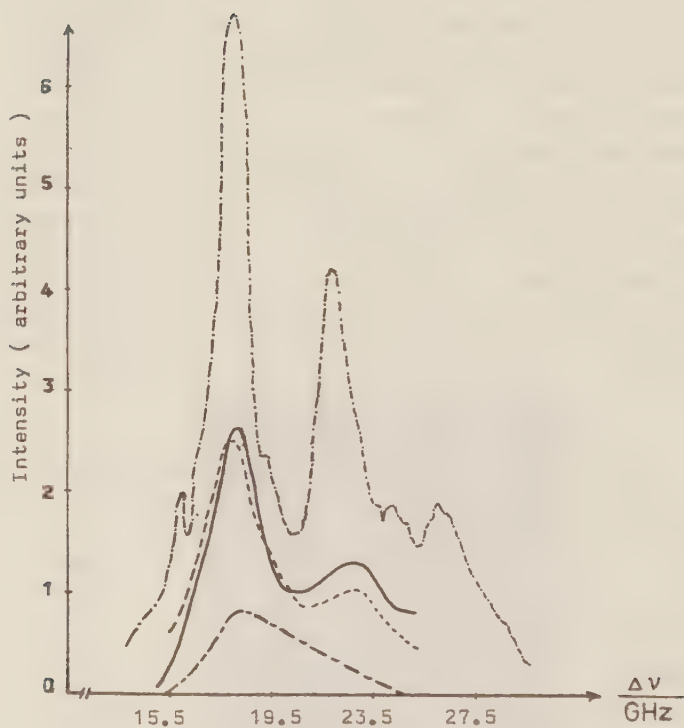


Fig.(4.8) Time resolved backscattered light spectra of three laser pulses in a shot. The spectrum (---) corresponding to the first laser pulse with SBS less than 10% is not splitted while the others with SBS of about 15% are splitted. The total time integrated spectrum (-.-) is also plotted for the purpose of comparison.

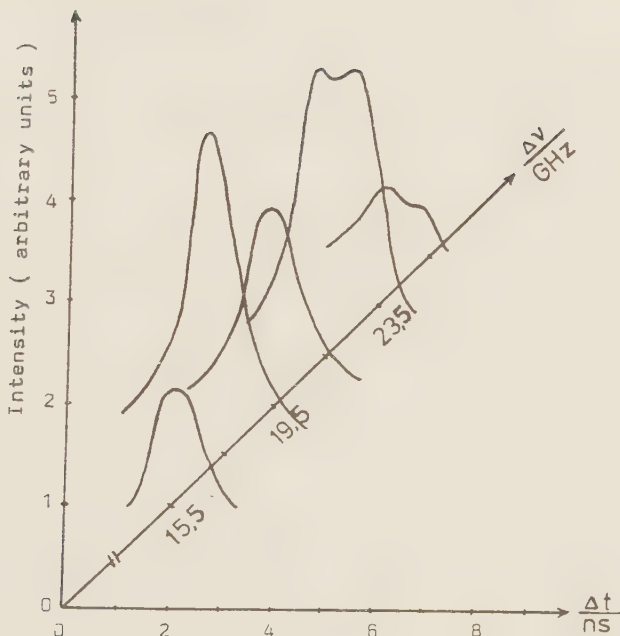


Fig.(4.9) A time resolved backscattered laser pulse spectra. SBS was about 12% and the spectra was a splitted one (shot # 39230).

- Rise of a non-splittd spectrum is normally quite sharp but the fall is quite slow, as shown in Fig.(4.10). It is in good agreement with the highly temporally resolved picture of the wavelength-integrated SBS laser pulse, Fig.(4.16). In this case, it was also observed that the rise of the pulse was quite steep but the fall was slow. The Fourier transform of the SBS pulse from the time domain to the frequency domain will yield a pulse of the form as it was observed in Fig.(4.10).

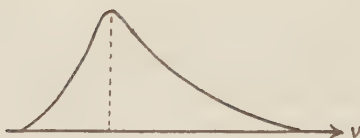


Fig.(4.10) A typical non-splitted spectra of a back scattered laser pulse (shot # 39117).

- As mentioned earlier in this section, the dependence of the frequency shift on the laser intensity will be now studied by analyzing a typical shot. Both the time integrated and time resolved spectra are shown in Fig.(4.11). Three different laser pulses were temporally resolved. All the laser pulses have different intensities and are indicated by I_i , I_j etc. producing a frequency shift of $\Delta\nu_i$, $\Delta\nu_j$, respectively. The plasma was stationary at the time of interaction for each pulse. Hence, it eliminates the possibility of the additional frequency shift due to Doppler effect. A study of large numbers of splitted and non-splitted spectra has shown that a better correlation between the laser intensity and the frequency shift can be established if it is assumed that the two peaks of splitted spectra are the sidebands about the main wavelength with a red shift of $\Delta\nu$. To check the validity of the $I_o^{1/3}$ and $I_o^{1/7.5}$ law, the frequency shift as shown in Fig.(4.11) has been plotted as a function of laser intensity in Fig.(4.12). It confirms our heating model by following $I_o^{1/7.5}$ law.
- Doppler effect was observed when the laser had interacted with the plasma in its dynamic phase. When the frequency shift corresponding to this Doppler effect was subtracted, the remaining red shift was that expected from the heating model and, in this way, proves the validity of our heating model also for the dynamic plasma phase.

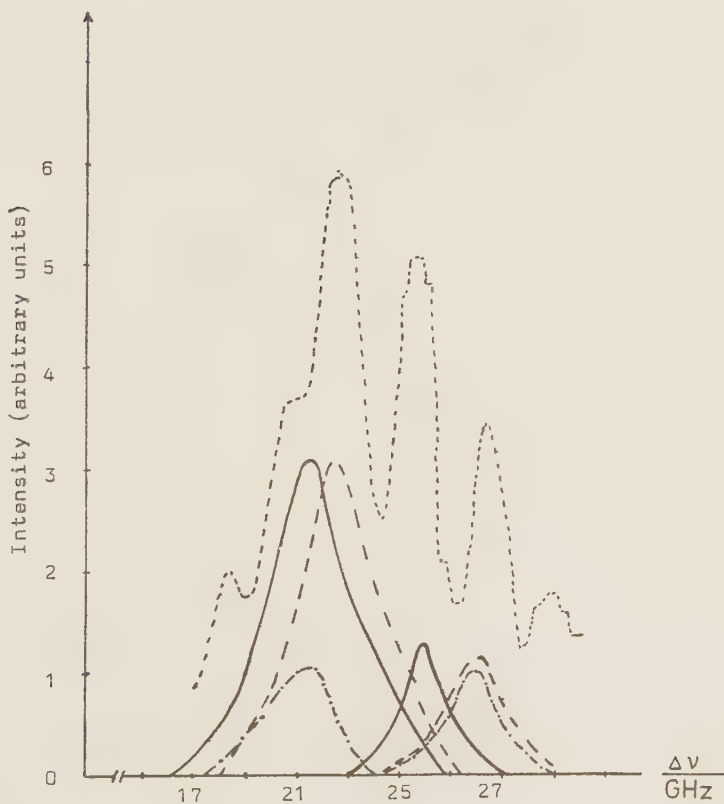


Fig.(4.11) Both the time integrated and time resolved spectra are plotted on the same scale. The time resolved spectra of three different laser pulses in a shot correspond to (—) 13% SBS, (---) 28% SBS, and (-·-) 17% SBS.

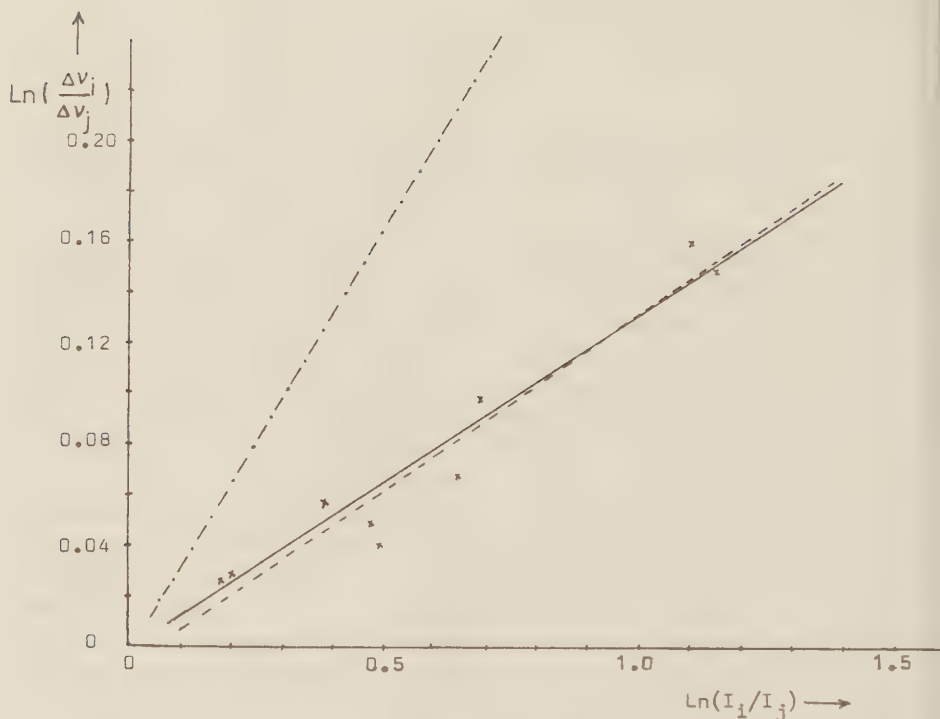


Fig.(4.12) The ratio between the frequency shift $\Delta v_i, \Delta v_j$ of the two back scattered pulses in several shots is plotted as a function of the ratio between the intensities of the corresponding laser pulses I_i, I_j . To check the dependence of the frequency shift on the laser intensity different possible exponent laws have been employed. A straight line (—) and a broken line (---) represent the frequency shift variation according to $I^{1/7.5}$ and $I^{1/3}$ laws, respectively. Crosses refer to the experimentally measured values. From the slope of the best line fitting (---) to the experiment data, the intensity dependence of the frequency shift was found to be $I^{1/7.2}$, which is in good agreement with the $I^{1/7.5}$ law.

Moreover, since this model proves to predict correctly the red shift dependence on intensity both below and above 10% back scattering, e.g. from where on an anomalous saturation mechanism sets in, it proves that this unknown process must be of a kind which does not interfere with the processes associated with plasma heating by the laser.

However, the splitting of spectra can not be explained by the heating model. In the hope to gain some insight into this unknown saturation mechanism the split spectra are dealt with, now.

4.4 Split Spectra

An important aspect of the spectrally resolved single pulse spectra which becomes evident for $\text{SBS} \geq 10\%$ is the appearance of a tail or second peak in the spectra within a frequency range from 4 to 5.5 GHz.

To explore the nature of the physical mechanism responsible for the splitting of the spectrum, splitted spectra were analyzed with highest temporal resolution. Both the total back scattered laser pulses and the corresponding wavelength channel signals were observed on a 7104 oscilloscope with 7A29 plug in unit at a 2 ns time scale. It limits the observation to only one back scattered laser pulse in a shot. The time resolution was better than 200 ps. A typical shot with an SBS pulse (signal upward) and the corresponding different wavelength channels (signals downward) are shown in Fig.(4.13).

The time differences which were deliberately added to avoid any overlapping between the different channels were measured by directly irradiating the array detector with the laser radiation. If this time difference is eliminated, then the signals overlap, Fig.(4.14) and the real time development of the different frequency channels in the plasma can be studied.

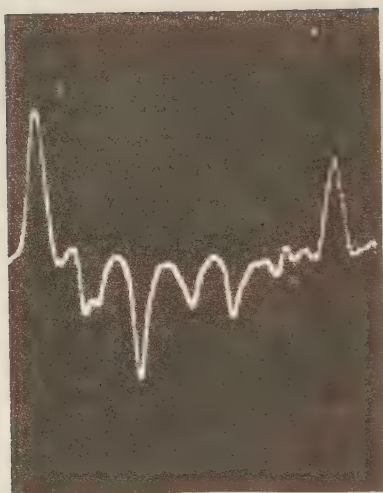
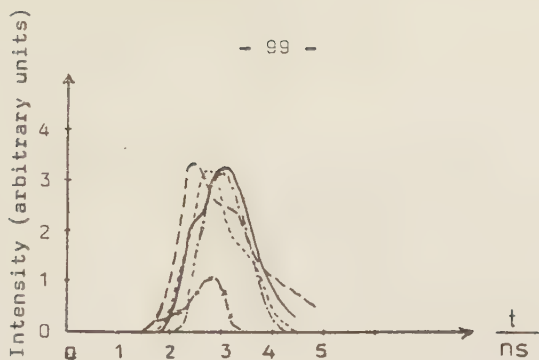


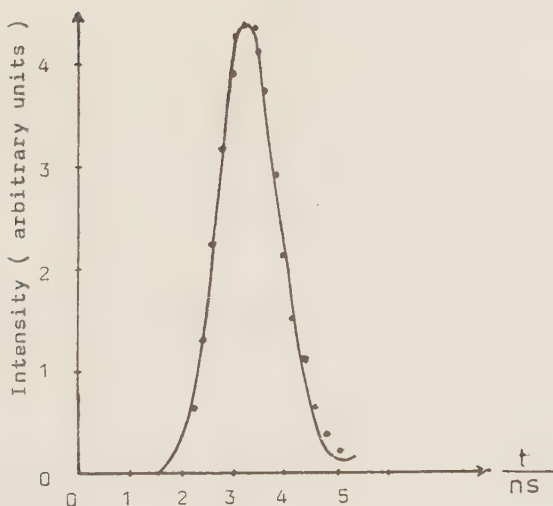
Fig.(4.13) Wavelength integrated backscattered laser pulse (signal upward) is spectrally resolved with the temporal resolution of better than 200 ps. The corresponding different wavelength channel signals are shown downward (shot # 39226),

Some important features can be immediately noted:

- signals at all the frequency channels start growing almost simultaneously, i.e. within ± 150 ps as shown in Fig.(4.14a)
- if all the different wavelength channel signals (Fig.4.14a) are added and the maximum is normalized to the peak of the corresponding backscattered laser pulse, then this pulse can be reproduced Fig.(4.14b). It proves the reliability of the measurements.



(a)



(b)

Fig.(4.14) (a) Time resolved five different wavelength channel signals corresponding to a backscattered laser pulse
 (b) summation of all the channel signals (dots) reproduces the wavelength integrated corresponding backscattered pulse (solid curve)
 (shot # 39212).

- another feature is a variation in the frequency difference ($\delta\nu$) between the two peaks of the spectrum. It was found that $\delta\nu$ decreases gradually as SBS increases. Such an effect can be seen in Fig.(4.15). It amounts to about 5.5 GHz, which is about 30% of the total frequency shift, near 10% SBS and it reduces to 3.9 GHz (about 20% of the total frequency shift) for 30% SBS. In other words, an increase in SBS of 20% reduces this frequency difference 28%. However, no correlation between $\delta\nu$ and Δt , i.e. the time difference between back scattering and the maximum compression of the plasma, was found.

There are two possibilities for the formation of the second peak in a spectrum. One possibility is the existence of two particle groups in the plasma with different velocities causing different Doppler shifts. It can be easily ruled out on the ground that no correlation between the spectrum splitting and Δt could be established, even when Δt was very long. Moreover, splitted spectra should always be observed irrespective of the amount of SBS. It is contrary to our observations.

Another possibility is the presence of any saturation mechanism responsible for the spectrum splitting and for a sharp bend in the SBS curve (Fig.4.1).

All the saturation mechanisms which can play some role in our experiment have been described in section (2.2). It has been shown, that, among them, only three saturation mechanisms, ion wave harmonics, rescattering and ion-wave trapping may be effective near 10% SBS in our experiment. Their possibility will be now checked in the light of the experimental results described in this Chapter.

Generation of the ion acoustic wave harmonics may be an effective saturation mechanism but it is unable to explain the splitting of the spectrum. Moreover, it is not possible to check the presence of ion wave harmonics in our backward or in forward scattering experiment as obvious from the parallelogram construction of Fig.(2.4).

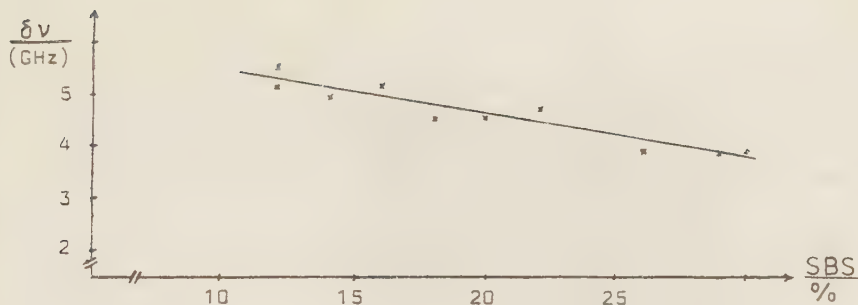


Fig.(4.15) Frequency difference ($\Delta\nu$) between the two peaks in a splitted spectrum as a function of the back scattering.

Besides ion wave harmonics, another possibility is the process of rescattering as large amounts of the back scattered laser light may produce rescattering in an underdense plasma of large extension. In an experiment described in sections (4.2.2) and (4.2.3), it could be demonstrated that this rescattering does not occur.

Besides the ion trapping possibility there exists a model of Forslund /8/ for the spectrum splitting. According to this model, the density modulation will be saturated when $\delta n/n \approx 1$ and a further increase in $\delta n/n$ will result in a frequency splitting of the standing wave. Due to pump depletion, the em-wave will propagate in a periodic structure as predicted by computer simulations. Both the incident and the back scattered waves will oscillate out of phase nearly sinusoidally at ω_0 but have a frequency splitting as

$$\delta\omega = \pm (\omega_{pe}^2 / 4\omega_0) (\delta n/n)$$

It becomes in our case

$$\delta\omega = \pm 6.8 \times 10^{12} (\delta n/n)$$

So, it predicts a very large amount of splitting and hence is unable to explain our experimental results. Moreover, the model is based on the assumption of pump depletion. But the ion trapping will occur before pump depletion as 10% SBS is not sufficient for pump depletion where the spectrum starts splitting. Hence, this model is not applicable to our case.

It is quite important from the SBS saturation point of view to check the existence of turbulence in the plasma. It increases the effective collision rate resulting in an increase in the ion temperature. Hence, the ion Landau damping becomes very effective. In this way the turbulent heating of the plasma leads to saturation of SBS.

To check the turbulence in the plasma before and after the laser interaction not only SBS was measured but also the backscattered pulses were resolved temporally.

A typical backscattered pulse recorded by a fast Tektronix 7104 oscilloscope and a fast pyroelectric detector UF₁, is shown in Fig.(4.16). Backscattered pulses were found to be sharper than the incident laser pulses. All the experiments reported to date had a time resolution of a few nanoseconds and hence were unable to realize the sharpness and the time difference between the incident and the backscattered pulses, as shown in Fig.(4.17).

From the pictures Fig.(4.16, 4.17) two important features of the backscattered pulses could be inferred to disprove the presence of turbulence.

It has been observed that the rise of the backscattered pulse was always faster than that of the incident laser pulse, while the fall follows the incident laser pulse which was almost gaussian.

Another feature of SBS could be inferred from the time coincidence between the incident and the backscattered pulses, as shown in Fig.(4.17). A time interval between exceeding SBS threshold and appearance of large backscattering was measured by a careful determination of the time of flight difference between the two pulses. It was found to be 300-400 ps, while the error was less than 100 ps. The observed time interval corresponds to SBS growth rates as they can be derived for a nonturbulent plasma in which SBS rises from thermal amplitudes by about nine orders of magnitude. Gellert / 13 /

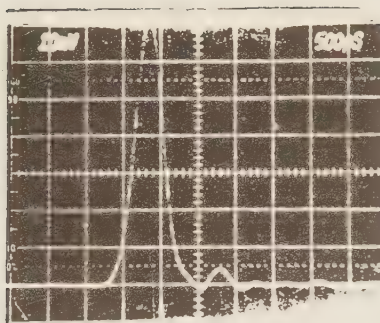


Fig.(4.16) A back scattered laser pulse as recorded by a fast pyroelectric detector UF1 and displayed on a Tektronix 7104 oscilloscope of about 350 ps rise time.

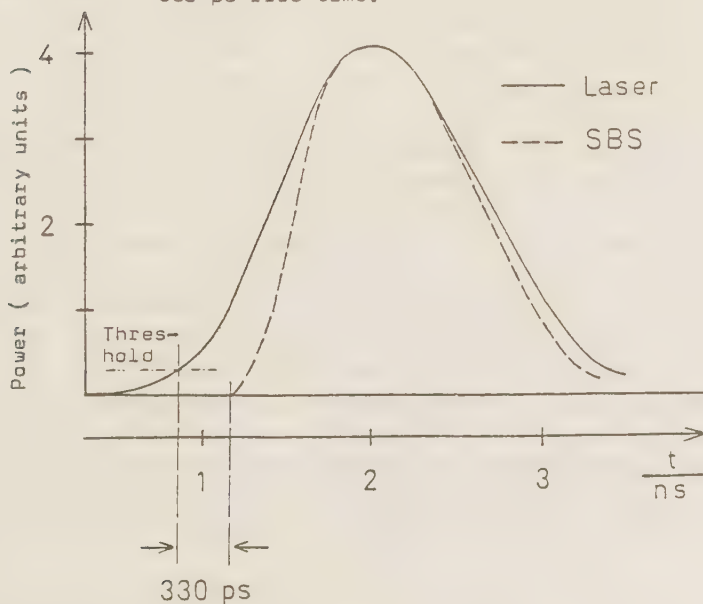


Fig.(4.17) Time coincidence between an incident and a corresponding back scattered laser pulse.

has shown that SBS starts from thermal amplitudes and this is another evidence for the absence of turbulence in our plasma.

Turbulence can also be ruled out from the observation of the back scattered light spectrum which was not much broadened and had only a frequency shift of (17 ± 1) GHz near threshold and (23 ± 2) GHz near 100% SBS. It has been measured /13/ and estimated /31/ that, when the laser interacts with the plasma, then the ion temperature increases very slowly but electron temperature increases very rapidly. In the presence of the turbulence, the ion temperature is expected to increase also rapidly due to the frequent collisions with the electrons. In the case of SBS, the frequency shift is directly related with $(K_B T_e + 3K_B T_i)^{1/2}$. A small increase in the ion temperature will result in much increase in the frequency shift. It is contrary to our observations. Even in the case of large SBS, ion temperature was not found to increase much. Hence, it can be concluded that there was no turbulence level in the target plasma before the laser interaction nor was produced by the laser heating of the plasma interaction volume.

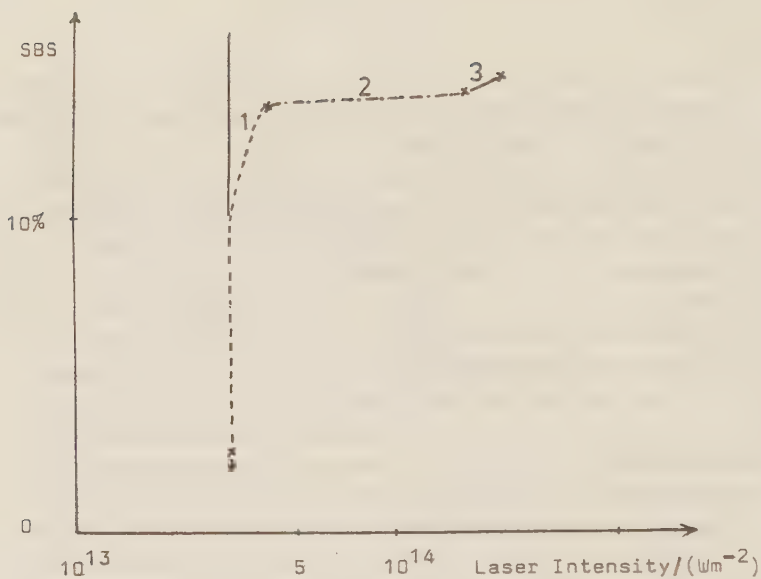
So, there is only one possibility left /73/, i.e., ion trapping and, next, an attempt will be made to associate the splitting of the spectrum with ion trapping.

As shown in Fig.(4.15), the frequency separation ($\delta\nu$) between the two peaks of the split spectra was found to decrease with SBS. On the other hand, Gellert's measurements /13/ show a tendency of decreasing density modulation with SBS. Hence, the decrease in $\delta\nu$ can be correlated with this density modulation $\delta n/n$ and, in turn, with ω_B (Fig.2.9) which is proportional to it. This confirms the validity of the ion trapping model and it means that the generation of side bands ($\omega \pm \omega_B$) about the main frequency ω is due to the ions bouncing with ω_B in the potential trough of the wave.

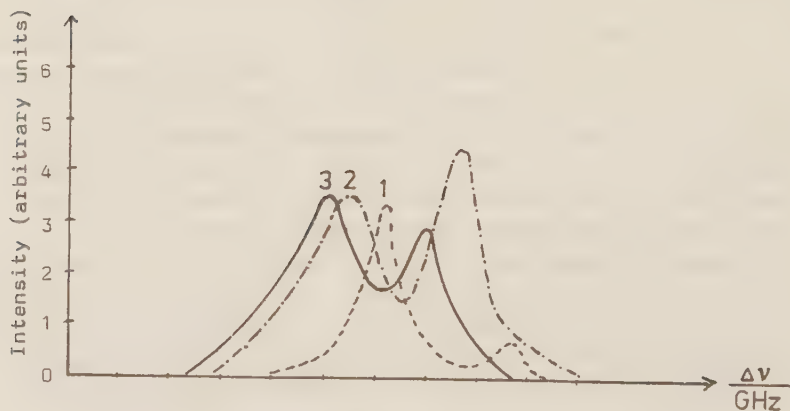
From Fig.(2.8), it is obvious that the density modulation will be about 10% near 10% SBS. In addition, it has been shown in section (2.2.5.1) that the density modulation of about 8% is required in our experiment for ion trapping to be effective. On the other hand, splitted spectra have been observed near 10% or more SBS. It implies that at the low amount of backscattering ($SBS < 10\%$) ion trapping is ineffective and the density modulation is less than 8%, which is the threshold for ion trapping. Hence, there was no enhanced damping. Near 10% SBS, the ion acoustic wave amplitude gets sufficiently large to trap a significant number of ions in its troughs resulting in the enhanced damping and in the generation of side bands.

The SBS light spectrum was found to contain both upper ($\omega + \omega_B$) and the lower ($\omega - \omega_B$) side bands. However, the amplitude of these bands was not found to be constant relative to each other and this aspect will be discussed now.

In Fig.(4.18), both the splitted spectrum and the position on the SBS curve are shown for three different laser pulses. It is quite interesting as in this case the amplitude of the upper side band changes to a large extent. However, it is easy to establish a correlation between the amplitude of the upper side band and the SBS curve. The SBS pulse 1 (Fig.4.18 b) is not much damped and its corresponding spectrum (Fig.4.18a) has a relatively small amplitude upper side band. On the other hand, SBS pulse 2 is highly damped and the amplitude of the upper side band is even higher than the lower side band. SBS pulse 3 is also damped but less than in case of pulse 2 and more than in case of pulse 1. Its corresponding spectrum has the amplitude of the upper side band comparable to the lower side band.



(a)



(b)

Fig.(4.18) (a) SBS as a function of laser intensity for three pulses within one and the same shot.

(b) Time resolved spectra of the same three pulses indicated by numbers 1,2 and 3. The horizontal scale is 1 mm = 0.2 GHz.

The ratio between the amplitude of the upper and lower side bands can be well explained by the trapped particle model based on Mima and Nishikawa theory as described in section (2.2.5.). On the basis of this criteria the amplitude of the upper side band with respect to the lower side band can be explained. According to their theory, the amplitude of the upper side band depends on the number of particles trapped which in turn depends on the amplitude of the ion wave. In case of pulse 2, much more ions were trapped and hence ΔN_t was greater than zero resulting in heavy damping, low amount of backscattering, and in the relatively large amplitude of the upper side band.

On the other hand, in case of pulse 1 only a small number of ions was trapped due to the small amplitude of the ion acoustic wave which implies that ΔN_t was less than zero resulting in the low amplitude of the upper side band.

It should be mentioned here that the gradual disappearance of the amplitude oscillations as observed / 38 / is due to those collisions by which the trapped ions are scattered. It may not be significant at low plasma densities at which the experiments /38,40,49/ on the IA-wave have been performed. As we have been working with high plasma densities, the effects of collisions too must be considered.

But it has been shown in section (2.2.5.3.) that collisions will not affect the ion trapping for a density modulation of 9% or more. This limit of the density modulation might be too low as in this estimate all types of collisions are taken into account. But all collisions are not effective for untrapping the particles.

The above considerations have shown that all the features of SRS can be well understood on the basis of the ion trapping model. A rough calculation has shown (section 2.2.5) that the acceleration of ions can absorb up to 50% of the

ion acoustic wave power resulting in the enhanced damping of the wave. Moreover, the spectrum splitting can be well understood on the basis of a simple model. In case of ion trapping, ions will be bouncing back and forth in the potential trough of the ion acoustic wave. It is quite similar to the case of an envelope function with some modulations superimposed on it, as depicted in Fig.(4.19a). A Fourier transform of such a time dependent function will yield according to modulation theorem a split spectrum, as shown in Fig.(4.19b). The two peaks of the spectra will be separated by twice the bounce frequency. It is in good agreement with our experimental observations.

From Gellert measurements /13/, the average value of the interaction length L as a function of SBS can be found. It helps to estimate the average density modulation per equation (2.46) for a particular value of SBS, as shown in Fig.(4.20a). It shows that $\delta n/n$ decreases sharply near 10% SBS and then later it becomes constant. Such a sharp drop in $\delta n/n$ can also be well understood in terms of the enhanced damping due to ion trapping. The larger the ion wave amplitude the more SBS will be produced resulting in the increase in the number of trapped ions. Hence, the enhanced damping will correspondingly increase which will damp the amplitude of the ion wave as obvious from the first part of Fig.(4.20a). But near 30% SBS a stage will reach when all the ions would have been trapped leaving no chance for further enhanced damping and hence $\delta n/n$ becomes constant. A further increase in the incident laser intensity and hence in SBS will increase the ion wave amplitude. As now there is no further enhanced damping so $\delta n/n$ will increase as can be seen in Fig.(4.20a).

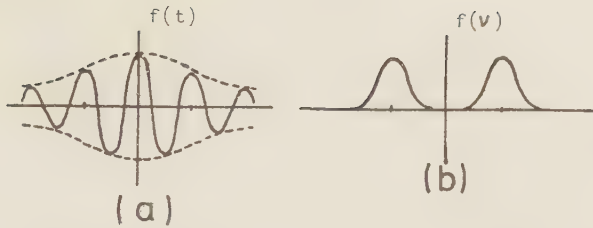


Fig.(4.19) (a) An envelope function (dotted curve) with a time period of $1/\nu_B$ representing the bouncing oscillations of the ions, superimposed by another oscillating function (solid curve) corresponding to an ion acoustic wave, and
(b) its Fourier transform, which is separated into two parts by twice the bounce frequency $(\nu_B) / 74\%$.

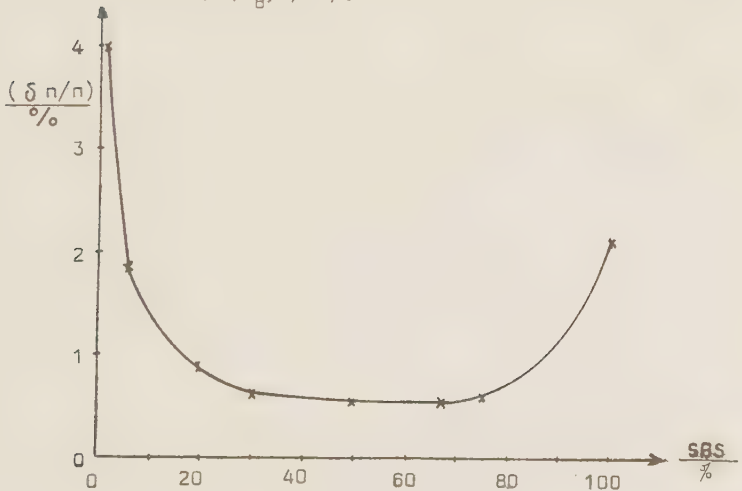


Fig.(4.20a) Density modulation as a function of SRS based on the measurements of the interaction length by Gellert /13 / which helps to estimate $\delta n/n$ by equation (2.46).

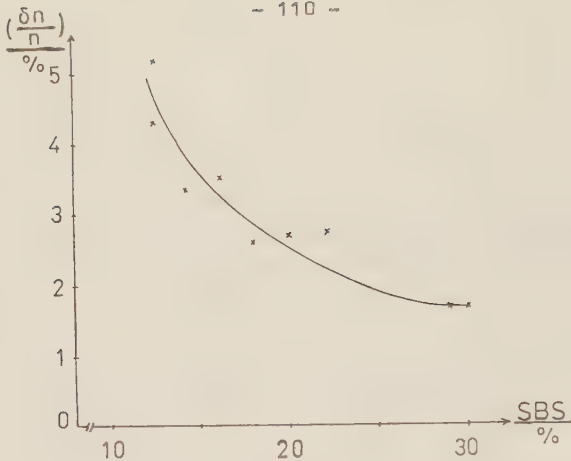


Fig.(4.20b) Density modulation as a function of SBS based on the measurements of δv as shown in Fig.(4.15).

Now our experimental results will be compared with Gellert's measurements. The knowledge of the frequency difference (δv) between the two peaks in a splitted spectrum is a direct measure of the bounce frequency. Hence, δv was measured and was found to decrease with an increase in SBS, as shown in Fig.(4.15). On the other hand, from Fig.(2.7), the electron temperature can be estimated corresponding to a particular laser intensity or SBS. The knowledge of bounce frequency from Fig.(4.15) and the electron temperature help to plot the density modulation per equation (2.58) as a function of SBS and is shown in Fig.(4.20b). It shows a decrease in the density modulation of about 60% as SBS increases from 10 to 30%. It is in good agreement with Fig.(4.20a). The assumptions made in the derivation of equation (2.58) can account for the difference between the absolute values of $\delta n/n$ in Fig.(4.20a) and Fig.(4.20b).

For ion trapping to be effective, the density modulation, initially in our case, should be about 10% (section 2.2.5.1). On the other hand, Gellert has measured a density modulation of less than 10%. It is in no way a contradiction as he has measured time average value of the density modulation which might be more than 10% initially but drops to a low value due to ion trapping.

Chapter 5

Comparison

Considerable experimental efforts have gone into studies of SRS. All the experiments can be categorized in two main groups, depending on the nature of the target, solid or gaseous. In case of a laser solid experiment the moving critical surface and the production of a magnetic field of several mega-gauss /75/ further complicates the problem. In such a case, the possibility of the stimulated scattering of laser light has been discussed elsewhere / 76 /. Therefore, the solid experiments will not be considered here. Even in the case of plasma, specifically designed as a model plasma, in one and the other aspects plasma parameters in the region in which SRS takes place were not well known and it makes difficult to interpret the results. For instance, Chen et al /77/ have observed the low level of back scattered radiation from an arc produced plasma whose parameters were measured before the laser interaction but at the time of interaction was only inferred.

On the other hand, Offenberger et al / 78 / have measured high level of SRS from a laser breakdown plasma formed in hydrogen gas. Again, no measurements of the plasma parameters were reported.

All the experimental results of SRS reported to date are very diverse in nature as shown in Fig. 5.1. To remove the discrepancies, other experiments will be compared with our experiment. A quantitative comparison is not possible as in other experiments plasma parameters were not known.

To the right in Fig.(5.1), it is not surprising to note a large back scattering in case of double laser pulse experiments on solid targets. In such cases, the interaction length was sufficiently large as measured interferometrically. On the other hand, some of the laser experiments on plasma targets also show large back scattering of up to 85% as observed at University of Alberta. But the back scatter from magnetically confined gas breakdown plasma at Universities of British Columbia and Washington and from arc pre-ionized plasma at University of California, Los Angeles is quite low.

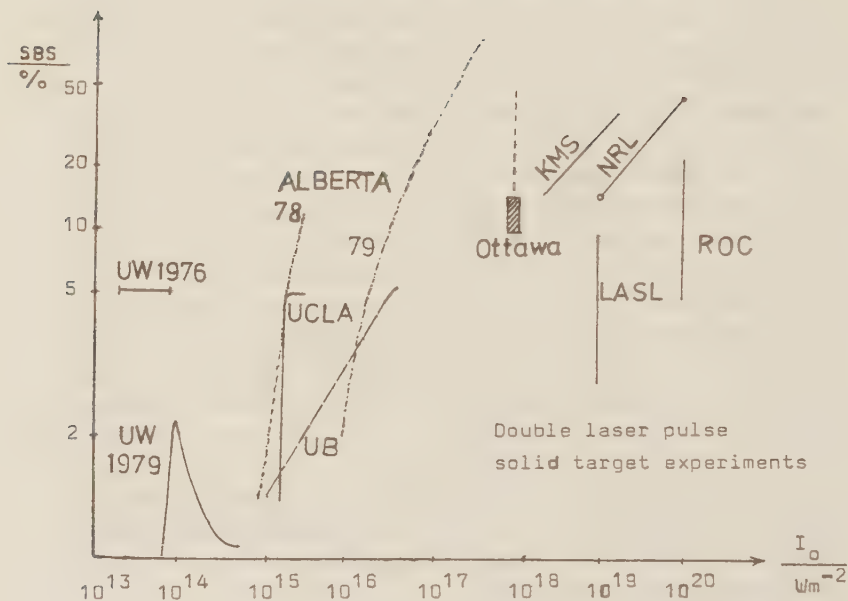


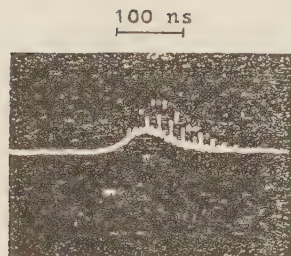
Fig.(5.1) Experimental measurements of back scattering versus laser intensity in case of gas and double laser pulse solid target as reported in the literature / 79 /, where

- UW = University of Washington
- UCLA = University of California, Los Angeles
- Alberta = University of Alberta
- UB = University of British Columbia
- Ottawa = National Research Council, Ottawa
- NRL = Naval Research Laboratory
- KMS = KMS Fusion, Inc., Michigan
- ROC = University of Rochester,
- LASL = Los Alamos Scientific Laboratory

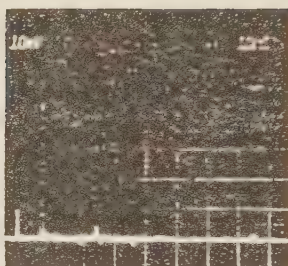
For the comparison of other experiments with our we should produce the same conditions. Hence, a laser plasma interaction experiment was performed when the laser was (1) injection mode locked, (Fig.3.7 c), (2) not injection mode locked and (3) mode locked by a cw laser. All the results reported in Chapter 4 correspond to the first case, while the case (2) in which the laser output is modulated and is irreproducible is shown in Fig.(5.2 a). In this case the FWHM of the laser was about 70 ns resulting in the ion heating of the plasma. As the rise of the laser pulse was quite slow, ions will get sufficient time to collide with the electrons resulting in low value of T_e/T_i . So, Landau damping will be effective and hence no back scattering was observed in the initial phase of the laser pulse. But laser has many sharp spikes (near its maximum). These sharp spikes will increase the electron temperature for a short time and T_e/T_i will become greater than unity and hence SBS threshold was achieved. So, some back scattering was observed as shown in Fig.(5.2 b). The results were not reproducible as the laser output was not reproducible. However, the amount of back scattering was always low. It has already been shown (section 2.2.8) that in such cases low amount of back scattering will be observed due to ion heating which is a very effective saturation mechanism and it quenches SBS immediately. Another spike again increases T_e/T_i and the process is repeated resulting in various back scattered pulses.

On the other hand, a smooth laser output (Fig.5.2 c) with even more power than in the previous case was not able to produce any back scattering. The reason is that in absence of spikes the ratio between electron and ion temperature (for long laser pulse) will never be much greater than unity. Hence, Landau damping will remain effective throughout the experiment and threshold will be very high for SBS. The possibility of SCS has already been ruled out due to the long e-folding time as described in section (4.1). So, no back scattering can be observed.

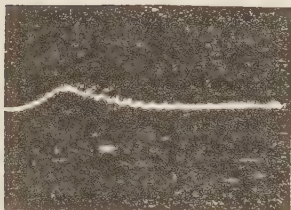
SBS threshold could only be achieved when $T_e/T_i \geq 3$. If this condition is satisfied then SBS will grow in the plasma till it saturates or the limit imposed by the Manley-Rowe relation is achieved. The growth of SBS not only depends on the laser intensity but also on the time when the saturation mechanism sets in the plasma.



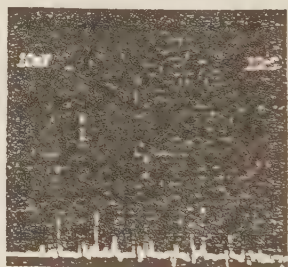
(a)



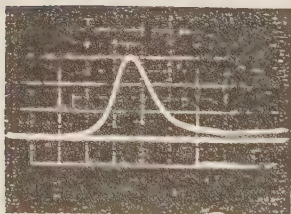
(b)



(a)



(b)



(c)

No SRS

Fig.(5.2) E-beam laser output (a) without IML and the (b) corresponding back scattered signals; E-beam laser output (c) with IML by a cw-laser. No back scattering was observed in this case.

Contrary to some other experiments, we have a train of laser pulses in a shot. It was assumed that time of 16.7 ns between the laser pulses is sufficient to restore the equilibrium conditions. Now it will be experimentally justified.

To find out the effect of the preceding laser pulse, two typical shots; one near threshold and the other with large SBS will be considered. In Fig. (5.3 a), the laser pulse was quite broad and had two peaks with a time difference of about 3 ns. The first peak in the laser pulse had produced a back-scattered pulse (SBS=3.3%), but the second could not produce any SBS. Again in the subsequent laser pulse, first peak had produced 3.8% SBS while the second had produced only 2.3% SBS. Similarly, in Fig. (5.3 b), all laser pulses had double peaks with an interseparation time of 2.2 ns resulting in two back-scattered pulses. The second peak had more intensity than the

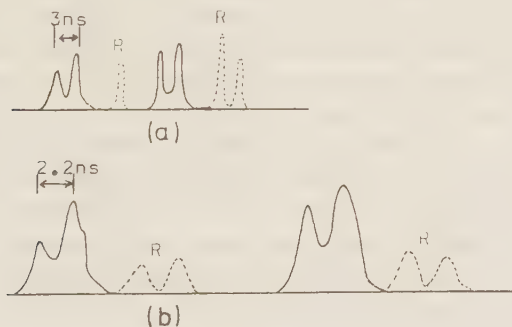


Fig.(5.3) Oscilloscope traces of incident and back scattered (R) laser pulses as recorded by a photon drag detector D_2 at two different time scales. Two different cases are shown here, one is (a) near threshold and the other is (b) large back scattering. The back scattered laser pulse in case (b) was much more attenuated than in case (a).

first pulse, inspite of that the first peak had produced more SBS (74%) than the second one which had produced a SBS of 62% only. This is contrary to the observations with single peak laser pulses (Fig.4.1), in which SBS increases with the intensity and is independent of the preceding laser pulse.

From these observations the following experimental facts can be established:

- a second laser pulse/peak preceded 2-3 ns by a comparable laser pulse/peak produces less backscattering,
- double peak laser pulse causes SBS to be double peaked,
- the time difference between the backscattered pulses was the same as between the two peaks in the corresponding incident laser pulses, and
- the backscattered pulses were somewhat shorter than the corresponding laser pulses of even more than 2 ns FWHM.

A simple explanation for the nature of the scattered pulse corresponding to a double laser pulse with a time separation of less than 3 ns is that the first peak of the laser increases the electron temperature rapidly but by the time when the second peak interacts with the plasma, ions have got sufficient time to gain some energy from the electrons due to collisions resulting in a decrease in T_e/T_i . The ion-Landau damping becomes effective and the growth of SBS will be reduced or completely stopped. With the rise of the second laser peak again T_e/T_i will rise and backscattering is observed. This time T_e/T_i will not increase much as ions were already heated, hence low SBS results.

As the interseparation time between the pulses was quite large, and up to 100% SBS was observed so it is justified to assume that in our case a preceding laser pulse does not affect the subsequent one. It was further confirmed by the low amount of frequency shift observed.

The specific objective in this thesis is to extract the information especially about the saturation mechanism contained in the backscattered light with emphasis on the time resolved spectroscopy. Not much work has been done on the temporal resolution of the backscattered light spectrum. To our knowledge only three reports exist for the case of laser produced plasma.

Ripin et al / 80 / have investigated a solid target with a Nd:glass laser. The maximum reflectivity was only 15%. They observed a complex structure in time and wavelength.

In another report, Turner and Goldman / 81 / investigated a solid target with a Nd:glass laser. They observed 9% backscattering. It was again not possible to understand the complex behaviour of the time resolved spectrum completely. The motion and inhomogeneities in the plasma and insufficient knowledge about the plasma parameters make, difficult to interpret the experimental results. But on the other hand, our time resolved spectroscopic study gives a very clear indication of the saturation mechanism present in our experiment and helps to understand our results. Even our model can be extended to explain the diverse nature of the results shown in Fig. (5.1).

Temporal modulation of the ion wave amplitude driven by SBS has been observed recently by Walsh and Baldis / 54 / reaching large amplitude at saturation of about 20%. The carbon plasma was produced by a Nd:YAG laser and was irradiated by a CO_2 laser. The plasma parameters were known before the interaction but not during the interaction. The Thomson scattering experiment with a Streak camera of 30 ps resolution shows a modulation of the ion wave with a bouncing frequency of 19 GHz. On the other hand, we have employed a fully ionized H_2 plasma and if we take into account the effects of all the relevant parameters, for instance, heating of the plasma by the laser, then our results are consistent with their results, as shown below.

In our case, the ratio between the ion wave frequency to the bounce frequency (v_{ia}/v_B) is about 6, and the electron temperature rises from 10 eV to 75 eV near 10% SBS. From the available data of Walsh and Baldis and by employing our heating model, it can be concluded that in their case electron temperature will rise from 100 to 600 eV and ion temperature from 100 to 150 eV when a 26 J, 1 ns CO_2 laser pulse interacts with a carbon plasma of $Z \approx 4$. So, the phase velocity of the ion acoustic wave in their case will be almost twice, hence in their case $v_{ia}/v_B \approx 2$. A difference of factor 3 between their results and ours can be well explained by the large bounce frequency in their case due to the high electron temperature (equation 2.58).

In addition to the amplitude modulation, they also observed the harmonic generation at large ion wave amplitude. Furthermore, it can be concluded from their results (Fig.2 in Ref. /54/) that on a time scale first the ion wave gets modulated and then the second harmonic appears. This confirms our model.

The experimental results can be divided in two groups. In one group, very low reflectivity (about 10%) was observed /80-84/. In such cases, normally laser pulse duration was quite large (about 40 ns) and it is quite possible that ion-saturation might have suppressed the instability. There are also a few reports in which about 50% back scattering has been observed /85-88/. In such cases, laser pulse duration was short (a few ns) and, ion saturation may not be effective but some other saturation mechanism might have competed with SBS or the interaction length was not sufficient enough for SBS to grow to 100%. It has been shown by Gellert /13/ that an interaction length of about 10 mm is required for 100% SBS. As sufficient data are not available, it is not possible to understand each experiment quantitatively. However, a

comparison of other experiments with ours helps to understand their diverse nature and the reasons which lead to 100% SBS in our experiment but not in other experiments. They are:

- large interaction length (SBS increases with the interaction length per equation (2.46))
- homogeneous plasma (inhomogeneous plasma has a high threshold)
- no hot spot in the laser output (otherwise filamentation may compete with SBS)
- short laser pulse (long laser pulse results in ion heating)

The above analysis indicates that a large amount of SBS will be produced in a large underdense plasma with small density gradients. It is of much importance as such conditions will be encountered in the corona of second generation, isentropic compression targets.

Chapter 6

Summary

To understand the widely different results reported in literature, a model laser plasma interaction experiment was performed to study the physics of scattering process under well defined conditions. In our experiment, a solid target was not employed as in the case of laser produced plasmas it is difficult to establish the conditions for the interaction due to the motion and inhomogeneities in the plasma.

For the basic understanding of the subject, both linear and nonlinear theories of the laser plasma interaction were briefly reviewed. It was assumed that the plasma is homogeneous and underdense. Among many parametric processes only Stimulated Brillouin Scattering (SBS), Stimulated Compton Scattering (SCS) and Stimulated Raman Scattering (SRS) were considered as they may scatter the laser radiation to a significant amount.

From the linear theory, threshold and the growth rate of SBS, SCS and SRS were determined which helped to identify the parametric process responsible for the laser light scattering. As SBS was found to be most significant among them, a further study of the subject was restricted to SBS only. It has also been shown that maximum scattering occurs in the backward direction, i.e. in the opposite direction to $\vec{k}_0$.

On the other hand, the nonlinear theory predicts the most significant feature of a scattering process which is its ability to scatter a certain amount of the incident energy. It was found that the plasma interaction length, the plasma density, and the ion density modulation level determine the amount of back scattered energy.

An attempt to explain our experimental results on the basis of the existing theoretical models did not succeed. It was found that these models are not of universal nature.

Although a large amount of SBS has been observed, yet it does not imply a complete absence of all the saturation mechanisms which may affect the growth of SBS drastically. Later the possibility of the saturation mechanism was confirmed experimentally. It was found that the amount of back scattering does not increase so rapidly with the laser intensity as it is expected on the basis of the nonlinear theory. A theoretical study of the saturation mechanisms has shown that only rescattering, harmonic generation, and ion trapping may be effective in our experiment. As harmonic generation has already been analyzed by Gellert /13/ and the existence of rescattering could not be experimentally established, ion trapping was discussed in detail. The density modulation as a function of T_e/T_i was determined while T_e/T_i , in turn depends on the laser intensity etc. and can be explicitly calculated for a particular laser intensity by the Heating Model of Landke /31/. The knowledge of the density modulation was necessary to estimate the bounce frequency of the ions. The threshold condition for ion trapping which depends on the bounce frequency, was established. It was found that in our experiment a minimum density modulation of 8% is required for the ion trapping to be effective. It was also indicated in chapter 2 that ion trapping will result in the enhanced damping and in the generation of the side bands due to the interaction between the ion-acoustic wave and the trapped particles bouncing in its potential trough. In addition to it, the effect of the collisions has also been considered as our plasma was not collisionless. Mima and Nishikawa theory /56/ for the electron trapping was extended to the case of ion trapping by including the ion dynamics. In this analysis all types of collisions have been taken into account. A rough estimate has shown that in our case collisions will become effective for a density modulation of 9% or less.

The above theoretical description was essential for the understanding of our experimental results. On the experimental side, a high power CO_2 laser system was developed with certain specific requirements. A clean laser plasma interaction experiment required the laser output to have a high contrast ratio, single longitudinal and transverse mode, and small beam divergence. Moreover, the laser pulse should not be of more than 2 ns duration to avoid any density variation in the interaction volume of the plasma during the experiment. In addition to it, the laser pulse should not be modulated, otherwise the results will not be reproducible. To avoid any filamentation instability which may compete with SRS, the laser output should be free of hot spots. Therefore, a lot of efforts were invested to obtain an unstructured focal distribution. On the other hand, sharp laser pulses were achieved by injection mode locking (IML) as described in detail in the Appendix. As a master oscillator, a helical TEA CO_2 laser of semi-confocal type was employed. A single longitudinal and transverse mode operation was achieved by introducing an iris inside the resonator and by replacing the front mirror with a resonance reflector. Moreover, a pulse cutting system was employed to cut a 200 ns laser pulse to a short pulse of a few nanosecond duration.

For injection mode locking, the output of the maser oscillator was injected into the slave oscillator. The slave oscillator was an e-beam CO_2 laser system with a resonator of the unstable configuration. The discharge chamber was filled with a gaseous mixture of He and CO_2 at a pressure of 2×10^5 Pa and the active volume was about 10 litres. N_2 gas was deliberately avoided in order to reduce the energy of the delayed pulses which may damage optical components.

The e-beam laser system was optimized and the characteristics of its output were studied. To check the experimental results, two different computer programs, as described in Appendix, were employed to predict the time evolution of laser output power, upper laser level and population inversion density, ambient temperature, and the pulse amplification in our e-beam CO_2 laser system, i.e. IML. The dynamics of the excitation and de-excitation processes in the gas mixture, the properties of the unstable resonator configuration, and the processes of the amplification of a few nanosecond laser pulse while propagating through a prepumped gas medium were considered. It was essential for the development of the computer programs. The computer model 1 predicts the e-beam laser output characteristics and helps in the optimization of the laser. In addition to it, this computer model furnishes the information about the time development of the upper and lower laser level which was essential for the description of IML by another computer model. It has been shown in Appendix that the computer results are in good agreement with the experimental results. It confirms that the laser was operating at its optimum performance and at the same time its output was satisfying all the requirements of a model laser plasma interaction experiment. Our laser system emits a train of nanosecond pulses with an interseparation time of 16.7 ns, which is equal to the round trip time of the resonator. In LPI experiment, a train of pulses was found to be advantageous in the sense that it has yielded several measurements in one shot and the time difference between the pulses was sufficient to restore the equilibrium conditions in the plasma.

The plasma was produced in a dynamical Z-pinch device with a total discharge energy of 20 KJ. The plasma density in our experiment was kept quite underdense although there was a provision to produce a plasma even above the critical density. The plasma was fully diagnosed by Finken et al / 69-71 /. It has provided significant information about

the plasma parameters required for the understanding of a LPI experiment. Most of the experiment was performed at a plasma density of $0.16 n_{cr}$ since an underdense plasma was desired and the plasma parameters were known for these conditions. To avoid the dynamical effects, laser was so synchronized with the plasma that the interaction occurs about the time of plasma maximum compression. In this way, some results of Finken et al were confirmed.

An homogeneous fully ionized hydrogen plasma of $0.16 n_{cr}$ was irradiated with a laser of maximum intensity of 10^{17} Wm^{-2} . The nature of the scattered radiation was identified as SBS. An attempt to observe SRS at a low plasma density of $0.08 n_{cr}$ has failed. The absence of both SRS and SCS was interpreted as a consequence of the short laser pulses.

SBS radiation both in the forward and in the backward direction was investigated in detail. The experiment was performed with a high temporal and spectral resolution of better than 200 ps and 0.6 GHz, respectively. The detection limit was 0.003 of the incident light and IRSA system was an order of magnitude more sensitive. In spite of that no rescattered radiation could be detected either in the forward or in the backward direction.

For the first time, back scattering power amounting up to 100% was observed. The threshold laser intensity was found to be in good agreement with SBS theory. The back scattered radiation was studied both temporally and spectrally. It has proved the absence of turbulence in the initial target plasma even during the laser heating.

A study of large number of time integrated back scattered light spectrum has revealed a red frequency shift of about 17 GHz relative to the incident laser radiation. It was found that the shift and width of the scattered light spectrum increases with SBS, which in turn depends on the laser intensity. A close study of the spectrum has confirmed that scattering is due to some parametric processes rather than due to the bulk motion of the plasma.

The frequency shift has permitted to determine the weighted sum of electron and ion temperature which has enabled Handke / 31 / to develop a heating model. This model determines explicitly T_e and T_i and was later experimentally confirmed by Gellert / 13 /.

Besides the backward scattering, forward scattering was also studied. It was found that no real forward scattering occurs rather than the frequency shifted light in the forward direction is the back scattered SBS signal travelling back to the e-beam resonator, being reflected there and again irradiating the plasma when it becomes transparent.

Inspite of all these valuable information, it was not possible to understand the large deviation of the back scattering curve, which becomes apparent near 10% SBS, from the theoretical one. Another paradox was the dependence of the frequency shift on the laser intensity. Handke heating model predicts an increase in the frequency shift with the laser intensity as $I_0^{1/7.5} \propto \Delta V$ while the theory of Forslund / 8 / predicts $I_0^{1/3} \propto \Delta V$ in case of strong coupling. On the other hand, as SBS increases the spectra get broadened and it impedes to associate the laser pulses with the peak of the back scattered light spectra within a shot. Thus making difficult to check the validity of $I_0^{1/3}$ or $I_0^{1/7.5}$ law in a time integrated back scattered light spectrum.

To explore the nature of the saturation mechanism causing anomalous damping near 10% SBS and to check the dependence of the frequency shift on the laser intensity, the back scattered light spectra were temporally resolved. It was found that SBS spectra split into two parts near 10% SBS from where an anomalous saturation sets in. This spectrum splitting (δV) corresponds to twice the ion bounce frequency. Moreover, δV was found to decrease gradually with the increase in SBS within a range of 10-30% SBS. It was also found that the frequency shift follows $I_0^{1/7.5}$ law.

All the possibilities which may cause the anomalous damping thus resulting in the spectrum splitting have been considered. It was found that all the experimental

results could be well understood on the basis of the ion trapping. It enables to develop a model for our experiment which is as follows:

The sharp laser pulse increases T_e/T_i much above unity quite rapidly making Landau damping ineffective and thus increasing the amplitude of the ion acoustic wave. As soon as the ion wave amplitude gets sufficiently large, a fraction of ions from the bulk of the distribution function can be trapped from 10% back scattering on, corresponding to an initial density modulation of about 10%. The trapped ions will get accelerated at the expense of the ion wave energy causing it to be heavily damped. In fact, at this point the back scattering curve shows the sharp bend. Due to the increased damping the ion wave will taper off and settle at a lower amplitude. Back scattering would also drop, were it not that the interaction length increases. From Gellert's work/13/, it is known, that, after this transition the density modulation settles at values well below 10% in this quasi-stationary phase with the interaction length increasing to produce the back scattering measured. In this subsequent quasi-stationary phase, the anomalous damping was found by Gellert to be determined by harmonic production setting in after sufficient ions have been trapped. Thus, for the final damping by ion wave harmonics to develop, this transitional phase of ion trapping must go before. Only during this initial phase, the ion wave has sufficient amplitude to feed bouncing oscillations strong enough to modulate the ion wave amplitude which in turn results in the splitting of the back scattering spectrum.

Also in other words, the results of this work complement the findings of Gellert. For the comparison with other experiments, the experiment was also performed when the laser was not injection mode locked and was mode locked by a cw laser. It has confirmed our model.

Our model not only explains our results but also helps to understand the very diverse results reported in the literature. It also points out the reasons for the observation of a large amount of SBS in our experiment but not in others. Moreover,

our results are in good agreement with the recent observations of Walsh and Baldis /54/ about the ion wave amplitude modulation and the second harmonic generation.

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APPENDIX

A.1 Specifications

(1) Z-Pinch Device:

A detailed description is given in Ref. / 69 - 71 /. Here, only some important characteristics will be given.

Electrode distance	50 cm
Material of the discharge chamber	quartz glass
Inner diameter of the vessel	18 cm
Total bank capacitance	40.5 μ F
Total inductance	30 nH
Discharge energy at 32KV	20.7 KJ
Maximum discharge current	1.2MA
Filling Pressure	33-665 Pa
Diameter of the dense plasma column at maximum compression	18mm
Time of maximum density and temperature	400 ns after initial gas breakdown; $k_B T_e \approx k_B T_i = 10 \text{ eV} \pm 10\%$

(2) Electron Beam Sustained CO₂ Laser:

Type: Maxwell e-beam laser system with large area cold cathode. Manufactured for KFA Jülich in 1973.

Cavity: Unstable resonator configuration

Beam Divergence: 0.06 mrad

Half Power Point (e^{-2}) : 100 μ m in the focal plane of a 0.7 m diffraction limited lens

Maximum Output Power : with injection mode locking 5 to 10 GW

Total Output Energy : 45 J

(3) Helical TEA CO₂ Laser:

Cavity : Semi-confocal type

Beam Divergence: 1.5 mrad

Output power $\leq$ 500KW (FWHM 200 ns)

FWHM : 2ns(typical) with PCS (Pulse Cutting System)

(4) CW CO₂ Laser :

Type:Sylvania Cw CO₂ laser Model 941

Beam divergence: 4.0 mrad

Beam diameter: 4 mm

Output power: 3W(typical)

(5) Energy Metre :

Type: GenTec Inc. Model LED-200-C

Sensitivity: 8.4V/J

Sensitive area: 3.8 cm²

Rise time : 5ms

Spectral sensitivity: 0.3 μ m- 30 μ m
range

Maximum energy density: 0.5 Jcm⁻²

(6) IRSA System:

Grating:

Firm: Bausch and Lomb

Size: 206 x 128 mm²

Serial Number: 1280-9-1-1

Catalog Number: 35-53-27-890

Blaze Angle : 36°52'

Free Spectral Range : 12 μm

Linear Dispersion : 24 Å/mm or 6.3 GHz/mm

Intensity of Maximum Ghost : 0.0025% of parent in the
first order of 0.55 μm

Lines / mm : 150 at 8 μm in first order

Mirrors:

1.5 m off-axis concave mirrors of Halley Inc.

Spectrometer Type:

1.5 m Czerny-Turner

Infra-red Camera with Pyricon:

Firm: Hamamatsu / Thomson Companies

Target Diameter: 16 mm

Spectral Range: 3 - 16 μm

Irsa System Resolution Limit: 0.6 GHz with 1:3 magnifi-
cation optics before camera

(7) Optics:

All infrared optics was made of NaCl. It has an absorption coefficient of $\alpha = 0.0013 \text{ cm}^{-1}$ and refractive index $\zeta = 1.49$ at $\lambda = 10.6 \mu\text{m}$. High power Copper and Nickel mirrors of Spawr Optical Research Inc. (USA) and some gold coated mirrors were used.

(8) Oscilloscopes:

Tektronix fast (1 GHz) oscilloscope 7104 with 7A29 plug in unit was employed for high temporal resolution. In addition to it, Tektronix storage oscilloscopes 7834 and dual beam oscilloscope 7644 were also used.

(9) Detectors:

Various photon drag detectors; PD-7410, PD-7455 of Rofin Co. and PD-30; pyroelectric detectors P-500, P-501 of Molelectron Co. and UF1 of Pulse Systems Inc. were employed to measure the incident and scattered laser power etc. An array detector of Pulse Systems Inc. was also used for the temporal resolution of the back scattered laser light spectrum. It consists of six UF1 detectors in a row. Each element of this array has nearly the same operating characteristics as the UF1 detector.

	PD-7410	PD-30	PD-7455	P-500	P-501	UF1
Rise time (ns) (10% to 90%)	0.4	1	1	0.4	0.33	0.05
Sensitivity at 10.6 μm (V/MW) with 50 Ω cable	0.2	0.03	0.08	10	17	5
Sensitive Area (mm^2)	16	25	200	1.2	0.8	1
Damage Threshold (MW/mm ² for 1 ns pulse)	0.5	-	1	20	0.2	0.1

A.2 Basic CO₂ Laser Theory

The development of a high energy CO₂ laser requires the arc-free excitation of large volume of gas mixtures at high pressures. Normally, a glow discharge at atmospheric pressure is unstable and rapidly develops into an arc. However, the arc formation can be inhibited by a separate volumetric pre-excitation of the discharge volume. In this approach, a stable discharge is produced in a large volume, high pressure gas mixture by the injection of a high energy electron beam /A1/ or u.v.-beam /A2/ to produce and maintain the necessary ionization to carry the current. For the optimal design of a single line high power pulsed CO₂ laser, the following aspects should be considered.

As regards the laser cavity configuration it is difficult to obtain lowest order gaussian mode from a stable resonator for a high energy and a large active volume laser. In such cases, unstable resonators are the best alternatives. It will be discussed in section (A.4).

If a high density of electronic charge within the laser medium before and during the application of the main pulse can be produced then a large quantity of the electrical energy into the laser will be deposited. First of all, Fenstermacher et al /A1/ showed that this process could be enhanced and the range of operation of high pressure CO₂ lasers be extended by direct deposition of electrons from an external source into the laser mixture. When the preionization electron density is sufficiently large, the main discharge may be operated at a much lower field strength than that required to operate a self-sustained discharge. In this case, the electron beam controls the main discharge and it is known as an electron-beam-controlled laser, or more simply e-beam laser.

Many detailed analyses /A1, A3-A7/ of the molecular gas lasers excited by means of high current e-beams and operating at high pressures have been carried out. All these studies led to the conclusion that an e-beam controlled configuration lends itself to scaling to large volumes and to high pressures and gives a high gain and high energy output as compared to all other excitation schemes. Hence, an e-beam excitation scheme was employed in our laser system (section A.5).

Pulsed atmospheric pressure CO_2 lasers possess a strongly collision-broadened gain bandwidth of the order of ~ 4 GHz. So, the device tends to oscillate on many longitudinal modes and simultaneously quite often on a number of transverse modes as well. For our experiment single mode operation was desirable.

There are various approaches to obtain single modes operation from a TEA (transverse electrical discharge at atmospheric pressure) CO_2 laser. Single longitudinal mode operation could be achieved by placing a cell containing SF_6 inside the resonator, or diffraction grating, or by the use of a fabry-perot interferometer. These mode-locking techniques are quite sensitive to adjustments and the tuning range is limited, for example, to the bandwidth of SF_6 or other absorbing gases. Moreover, many aspects of the passive mode-locking are not well understood and it could not be extended to high energy lasers.

The best method to control the output of a high power laser is injection mode locking (IML). It is achieved by injecting a pulse from a small, low power, well stabilized oscillator into a high power laser, generally called a "slave oscillator". The effect of the injected pulse is to "lock in" and stabilize the frequency of the slave oscillator. First of all, Basov et al /A8/ locked the output frequency of a ruby laser by another one. During last few years, IML has been the subject of great interest /A9-A11/. It has been reviewed by Corkum /A11/ and he found IML to be the most suitable technique for obtaining the required output from a variety of lasers, with apertures up to 20 cm notably for the CO_2 , Nd:YAG, and Dye lasers. So, an injection mode locked laser system was developed, here, for the study of the parametric processes in the plasma.

For an optimal design of a laser, with certain specific features (section 3.1), numerical modelling is essential. It will be briefly described in the next section.

A.3 Numerical Modelling of the CO₂ Laser

Now-a-days the number of laser parameters has increased to such an extent that it is almost impossible to make a laboratory search for an optimal design of a laser. There has been great interest in modelling the behaviour of CO₂ lasers as it plays an increasingly important role in the design of a laser. With the help of these models, it is possible to obtain valuable information about the laser which is essential for an optimal design of a laser with the required features.

Originally, Witteman /A12/ presented a numerical model in 1966 to predict the gain in low pressure cw laser. Manes and Sequin /A13/ modified this model for TEA-operation by introducing a cavity intensity equation and pumping terms. This five equation kinetic model describes the temporal changes in the vibrational modes of CO₂ and N₂, and the fifth equation is for the cavity field intensity. Later, Davies et al /A14/ modified this model. In their model, three more equations were included to describe the CO vibrational mode, to allow variation in the temperature and to include an equation to describe the evolution of the population of the 100-state of CO₂.

"Tlaser"-and "pulsamp" computer programs of Davies et al /A14/ and Roberts et al /A15/, were modified here to predict the time evolution of output power, upper laser level and population inversion density, ambient temperature, and the pulse amplification in our e-beam CO₂ laser system, i.e. IML.

For the development of a computer program of a CO₂ laser, some insight into the dynamics of the excitation and de-excitation processes in the gas mixture is required. So, here the population distribution in different levels and the physical mechanism of an e-beam CO₂ laser will be briefly reviewed.

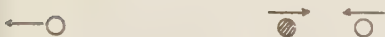
Generally, molecules are more complex than atoms but CO₂ is an unreactive, relatively simple polyatomic molecule as compared to other molecules. Its energy levels originate from electronic motions, vibrational motions and rotational motions. In general, in a CO₂ molecule not only the electrons will be excited but the atoms can also vibrate with respect to each other and, moreover, the molecule as a whole can rotate, as illustrated in Fig.(A-1). The CO₂ molecule has three modes



Symmetric Stretching Mode



Two Unsymymmetric Bending Modes



Asymmetric Stretching Mode

Fig.(A-1) Normal vibration of CO_2 , shaded circles represent carbon atom and blank circles refer to oxygen atoms.

of vibration, each giving rise to a series of energy levels: these are the symmetric stretch, the bending, and the asymmetric stretch modes. On the other hand, electrons have widely separated energy states and may give laser action in uv and in visible region. Atoms in CO_2 molecule can vibrate about their equilibrium position in such a way that electronic levels will not be perturbed. Vibrational frequencies are in the near and in the infra-red region. Hence, a CO_2 molecule may have many vibrational-rotational energy levels not found in any single atom. Vibrational levels are more closely spaced than electronic levels and rotational levels are still more closely spaced than the vibrational levels. The vibrational levels of CO_2 relevant to laser action are shown in Fig.(A-2)

In a CO_2 Laser, the molecule is initially pumped from its ground level to higher electronic and vibrational-rotational levels. It trickles down by relaxation through different levels, as shown in Fig.(A-2), depending on the relaxation rates and life-times of the levels.

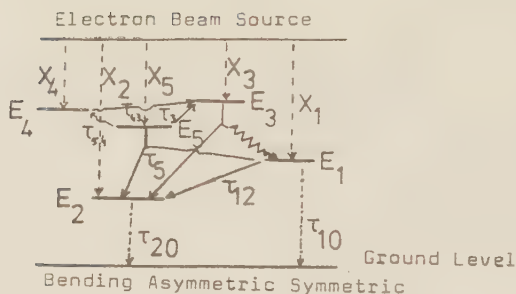


Fig.(A-2) Energy level diagram of the N_2 , He, and CO_2 molecule. X's are the excitation rates and τ 's are the relaxation times of the levels. For the sake of simplicity, the rotational levels are not shown here. Solid lines represent vibrational energy exchange between modes, (---) dotted lines refer to the relaxation of vibrationally excited molecules to the ground state, dashed lines refer to the electron excitation of vibrational modes, and wavy line represents the laser transition. E_1 , E_2 , E_3 are respectively the symmetric, bending, and asymmetric modes of CO_2 , and E_4 and E_5 are the vibrational modes of N_2 and He.

Molecules other than CO_2 may also help in lasing, for instance, He and N_2 . These molecules may also be easily excited by electron impact excitation and when they collide with CO_2 molecules, they give up their energy to them by resonant energy transfer process. Hence, they provide an efficient pumping mechanism. The vibrational state of a molecule may be described by the spectroscopic notation (n_1, n_2, n_3) , where $n_{1,2,3}$ are the number of quanta of frequency $\nu_{1,2,3}$, respectively, and 1 is the degenerative vibration ($1=n_2, n_2-2, \dots, 1, 0$).

Transitions occur between pairs of rotation-vibration levels according to well defined selection rules giving rise to

many possible transitions between $9\text{ }\mu\text{m}$ and $11\text{ }\mu\text{m}$. These selection rules are for:

vibrational transitions: $\Delta n=1$, $\Delta l=0, \pm 1$

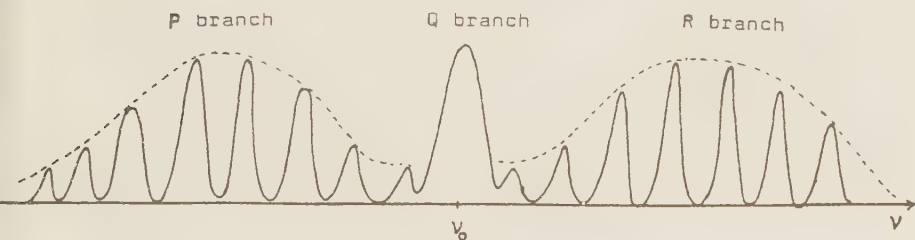
rotational transitions: $\Delta J=0, \pm 1$,

where J' is the rotational quantum number.

Vibrational and rotational transitions are characterized by the change in J as

$P (\Delta J' = -1)$, $Q (\Delta J' = 0)$, $R (\Delta J' = +1)$

and individual lines are denoted by $P(J')$, $Q(J')$, and $R(J')$, as shown in Fig.(A-3). It shows the transition between the two vibrational levels with account taken of the rotational splitting. The transition which in the absence of the rotational energy would have consisted of a single line at ν_0 actually



g.(A-3) Possible transitions between two vibrational levels, which in the absence of rotations would consist of a single line at ν_0 , and it will be made up of two sets of lines, P-branch and R-branch transitions. Q-branch transitions are not allowed, here, and hence will not contribute.

has two groups of lines, P- and R-branch. CO_2 is an exactly symmetrical molecule and has no nuclear spin, so $\Delta J=0$ (or Q-branch) transitions are not allowed. Moreover, J' to $J_{\pm 1}'$ transitions within a single vibrational state are not allowed, as mentioned above. So, here we dealt with a laser in which oscillation takes place between two vibrational levels in CO_2 .

After this brief review of the excitation processes in a CO_2 laser, now the characteristics of the e-beam resonator will be briefly described.

A.4 E-beam Resonator

To develop some physical understanding of the behaviour of our laser, it is useful to study first some of the properties of its resonator.

The e-beam laser system has an unstable resonator configuration with confocal mirrors. Unstable resonators can have large active volumes even in short resonators with high fresnel numbers. and are most suitable for high power laser systems. The theoretical description of the confocal unstable resonator is simpler relative to its counterpart in stable geometry and this configuration is unique in the sense that it radiates a parallel beam.

Resonators are described by four parameters, L' , R_1 , R_2 , and $\mathcal{G}$, where $\mathcal{G}$ represents the alignment of the mirrors. These geometrical parameters affect output coupling, far field intensity distribution, beam power and mode structure and will be discussed briefly. An optical resonator, as shown in Fig.(A-8), is generally characterized in terms of the curvature parameters as $/A16/$:

$$g_i = 1 - (L' / R_i) \quad (A1)$$

and the Fresnel number

$$N_{iF} = a_i^2 / (\lambda_0 L') \quad (A2)$$

where,

λ_0 = wavelength of the laser

$i = 1$ represents the rear mirror

$= 2$ represents the output coupling mirror

a_i = radius of the i th mirror

L = distance between the mirrors

R_i = radius of curvature of the i th mirror defined as positive if the centre of curvature lies towards the interior of the resonator.

Fresnel number is a dimensionless parameter describing the size of the mirrors of an optical resonator. Physically, it describes the number of Fresnel zones of a plane wave on one mirror of the resonator as observed from the centre of the other mirror.

A resonator was designed with the following parameters:

$a_1 = 50$ mm, $a_2 = 25$ mm, $R_1 = 10$ m, $R_2 = -5$ m

So, in our case,

$N_1 F = 94$, $g_1 = 0.75$, $g_2 = 1.5$ (A3)

Such a large Fresnel number indicates that diffraction losses will not be effective in discriminating between the various transverse modes of the cavity and a multimode output will emerge. However, a single mode operation can be achieved by reducing the Fresnel number to unity. It is possible for the cases where the output power is of not prime importance. This method was employed in the case of INL for master oscillator (section 3.1). But for a LPI experiment, a high power single mode laser is required. The only known possibility is the use of an unstable resonator / A17 /. An unstable resonator corresponds to a divergent periodic focusing system hence a lowest order mode expands in it on repeated bounces. Therefore, it is possible to achieve a single mode operation even for large mode volume since the laser output is taken as a diffraction coupled beam passing around rather than through the output coupling mirror.

The resonator will be stable if $0 \leq g_1 g_2 \leq 1$, otherwise unstable. The most useful form of unstable resonator is the confocal unstable resonator introduced by Krupke et al / A16 /. This configuration is normally preferred in high power lasers as it yields a collimated output beam. Unstable configuration is further categorized in two branches as:

$$\begin{array}{ll} g_1 g_2 \geq 1 & \text{positive branch} \\ g_1 g_2 \leq 0 & \text{negative branch} \end{array} \quad (A4)$$

In our case $g_1 g_2 \approx 1.1$. A positive branch configuration was selected as in this configuration light will not be focussed in the resonator which is undesirable in high power lasers.

In unstable resonators, spherical waves travel from one mirror to another but conventional Fresnel number describes a plane wave. So, Siegman /A17/ defined an effective Fresnel number based on the Fresnel zones of the actual spherical wave at one mirror observed from the centre of the other mirror. The effective Fresnel number is defined for the positive branch as:

$$N_{eq} = N(M-1)/(2M^2) \quad (A5)$$

where, M is the geometrical magnification factor and is the ratio between the radii of the curvature of the rear mirror to the front mirror. In our case, $M=2$. It is directly related with the geometric output coupling parameter δ as:

$$\delta = 1 - M^{-2} = 0.75 \quad (A6)$$

Substituting the values of N and M , the effective Fresnel number can be estimated for our case as:

$$N_{eq} \approx 11.7 \quad (A7)$$

The effective Fresnel number gives an indication for the resonator mode losses. The resonator mode losses will be maximum at integral values of the effective Fresnel number while single mode operation is favoured near the half-integral values of N_{eq} /A18/. But in our case, the equivalent Fresnel number is not a half-integer value. However, it has been shown /A19/ that for the output coupling greater than 60%, near diffraction limited, far field intensity distribution is insensitive to the particular value of equivalent

Fresnel number for large values of this parameter. Our selection of resonator parameters can be justified as in our case both the parameters, N_{eq} and δ , are quite large.

It is not easy to align the two mirrors of a resonator, perfectly. It is interesting to find out the range of the misalignment in which output does not deteriorate to a significant amount. It could be described in terms of the two angles, ϑ_a and ϑ_b . ϑ_a is the first critical angle, when the mirror is so much tilted that no power is coupled out of the coupling mirror.

ϑ_b is the second critical angle which occurs when the optic axis passes beyond the edge of one of the mirrors. For angles exceeding ϑ_b there exists no possibility of exciting modes characteristic of the aligned unstable resonator. These angles are defined for the positive branch configuration as /A16/:

$$\vartheta_a = \frac{2a_1}{L} \left(\frac{M-1}{2M} \right)^2 = 2.5 \text{ m rad (for our parameters)} \quad (A8)$$

$$\vartheta_b = \frac{a_1}{L} \left(\frac{M-1}{2M} \right)^2 = 1.25 \text{ m rad (" " ")}$$

It shows that our system was not much critical to alignment, and is advantageous for the stable laser output.

Although the output mode of the unstable resonator is a uniform symmetric ring in the near field, this transforms into a single main beam lobe in the far field with no hole in the axis. Intensity distribution of this far field resulting from circularly symmetric radiation emanating from an annular coupling mirror is given by /A20/

$$I(\psi) = \left| \pi (1 + \cos \psi) j \frac{\exp(-j2\pi L_R / \lambda)}{L_R} \int_a^M r f(r) J_0\left(\frac{2\pi r \sin \psi}{\lambda}\right) dr \right|^2 \quad (A9)$$

Here, the far field intensity distribution has been written in terms of an integral expression involving the near field intensity function $f(r)$, where

ψ = divergence angle

L_R = distance to the receiver located in the far field

r = radial variable at the near field source

$2a$ = inside diameter of the near field annulus (50 mm)

J_0 = zeroth-order Bessel function

$D = 2Ma$

The above relationship for the far-field intensity distribution has been computed and the result is shown in Fig. (A-4). In our

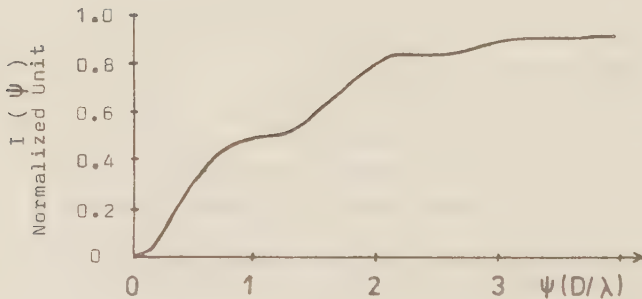


Fig.(A-4) The integrated far field intensity is plotted as a function of the divergence angle ψ in units of (D/λ) , /A20 /.

case; $\psi = 0.06$ mrad, $D = 100$ mm, $\lambda = 10.6 \mu\text{m}$, hence $\psi(D/\lambda) = 0.57$. It shows that almost 40 percent of the power is contained within the central lobe of the resonator for an unstable resonator with a geometric magnification of 2.0. It is helpful in the determination of the laser intensity in a LPI experiment.

After a brief description of the physical properties of the unstable resonator, the physical mechanism responsible for the laser action in an e-beam CO_2 laser will be considered. It will help to estimate the various important parameters which will be later required to solve the rate equations describing the kinetics of the CO_2 laser through a computer program.

A.5 Physical Mechanism of the E-beam Laser

In this subsection, the process of laser action will be briefly discussed and some parameters will be estimated for a computer program.

The discharge chamber of the e-beam laser has a gaseous mixture of He and CO_2 at a pressure of 2×10^{-5} Pa. The gaseous mixture resides in a steady electric field E_s of the sustainer bank and was irradiated by a high energy pulsed electron beam. Each primary electron creates a swarm of secondary electrons in the gas generating a highly uniform discharge even above the atmospheric pressure in a quite large volume of about 10 litres. The average energy of the secondary electrons is controlled by the magnitude of the imposed electric field chosen such that they efficiently pump the gas molecules into the vibrational levels.

The equation describing the production of electron-ion density is /A21/:

$$dN_e/dt = S - \alpha N_e N_i - \beta N_0 N_i \quad (A10)$$

where, N_e (N_i) is the electron (ion) density and N_0 is the neutral-gas density, S is an electron-ion source term produced by the injected high energy electrons ($\text{cm}^{-3} \text{s}^{-1}$), β is

the attachment coefficient (s^{-1}), α is the recombination coefficient ($\text{cm}^3 \text{s}^{-1}$). $\alpha N_e N_i$ represents the recombination loss per cm^3 per sec due to the existing densities of electrons and ions. The loss by attachment to neutral molecules is represented by $\beta N_0 N_e$. For the sake of simplicity, here it has been assumed that all molecules have the same recombination and attachment properties. Otherwise, a set of equations would be required to take into account the different properties of He and CO_2 . Charge neutrality condition allows to assume that $N_e \approx N_i$, so that

$$dN_e/dt = S - \alpha N_e^2 - \beta N_0 N_e \quad (\text{A11})$$

The solution of this equation is

$$N_e = -\frac{\beta N_0}{2\alpha} - \frac{(\beta^2 N_0^2 + 4\alpha S)^{1/2}}{2} \frac{\text{Tanh}\{-(t/2)(\beta^2 N_0^2 + 4\alpha S)^{1/2}\} - \beta N_0 / (\beta^2 N_0^2 + 4\alpha S)^{1/2}}{1 - \text{Tanh}\{-(t/2)(\beta^2 N_0^2 + 4\alpha S)^{1/2}\} \beta N_0 / (\beta^2 N_0^2 + 4\alpha S)^{1/2}} \quad (\text{A12})$$

which reduces to (for the steady case)

$$N_{e0} = -\frac{\beta N_0}{2\alpha} + \frac{(\beta^2 N_0^2 + 4\alpha S)^{1/2}}{2\alpha} \quad (\text{A13})$$

A dominant loss mechanism is the recombination of the electron-ion pairs, so the effect of the attachment could be neglected, / A1 /, i.e. $\beta/\alpha \longrightarrow 0$. Hence, (A12) and (A13) reduces to:

$$N_e = \left(\frac{S}{\alpha}\right)^{1/2} \text{Tanh} (S\alpha)^{1/2} t \quad (\text{A14})$$

$$N_e = (S/\alpha)^{1/2} \quad (\text{A15})$$

Equation (A11) can be solved for $\beta/\alpha \rightarrow 0$, if S and α are known. They can be determined by the following considerations: Normally, a 50 KV potential was applied at the anode in the discharge chamber and the anode-cathode spacing (distance between anode of the sustainer bank and the Ti foil) was about 10 cm (Fig.3.8). Hence,

$$E_s/N = 2 \times 10^{-16} \text{ V cm}^2 \quad (\text{A16})$$

at a gas density of $2.5 \times 10^{19} \text{ cm}^{-3}$ per atmospheric pressure. The drift velocity is the directed velocity of the electrons in the anode to cathode space of the discharge when they are exposed to an electric field, E_s of the sustainer bank, for the above value of E_s/N , and for a gas ratio

of He:CO₂ = 2.5:1, the electron drift velocity can be determined from the Fig.(A-5)

$$V_d = 5.8 \times 10^6 \text{ cm s}^{-1} \quad (\text{A17})$$

and the average electron energy can be estimated from Fig.(A-6) as:

$$\langle U \rangle = 1 \text{ eV} \quad (\text{A18})$$

The knowledge of $\langle U \rangle$ helps to estimate the recombination coefficients for CO₂ and He according to the following relations:

$$\alpha_{\text{CO}_2} = 1.7 \times 10^{-8} / \langle U \rangle \quad (\text{cm}^3 \text{ s}^{-1}) \text{ for CO}_2 \quad (\text{A19})$$

$$\alpha_{\text{He}} = 3.0 \times 10^{-7} / \langle U \rangle^{1.5} \quad (\text{cm}^3 \text{ s}^{-1}) \text{ for He}$$

The average value of α is:

$$\langle \alpha \rangle = 2.15 \times 10^{-7} \text{ cm}^3 \text{ s}^{-1} \quad (\text{A20})$$

Now, an attempt will be made to determine S. It can be evaluated if the secondary electron density N_{eo} is known, which requires in turn the knowledge of the electron beam current. A(-200 KV) pulse was applied to the cathode of e-beam gun by means of a pulsed transformer (Marx-generator) generating a plasma which expands towards the anode with a velocity of about 4 cm/ μ s /A23 /. The moving plasma can be considered as the virtual cathode and it can be assumed that the ensuing current flow is space-charge limited, and the average gap spacing is half-way between the e-beam gun cathode and the window of Ti-foil. As a first approximation, the total current density flow should follow Child-Langmuir law

$$J_p = 2.33 \times 10^{-6} V^{3/2} \left(\frac{d}{2} \right)^{-2} \quad \text{for planar cathodes} \\ = 11.7 \text{ A cm}^{-2} \quad (\text{A21})$$

where, V is the e-beam gun cathode voltage which was 250 KV and d is the gap distance between e-beam cathode and the window { which was about 100 mm in our case as described in section (3.11)}. The current density (J_d) entering in the discharge chamber depends on various losses and is given by

$$J_d = \alpha_1 \beta_1 \gamma_1 J_p \quad (\text{A22})$$

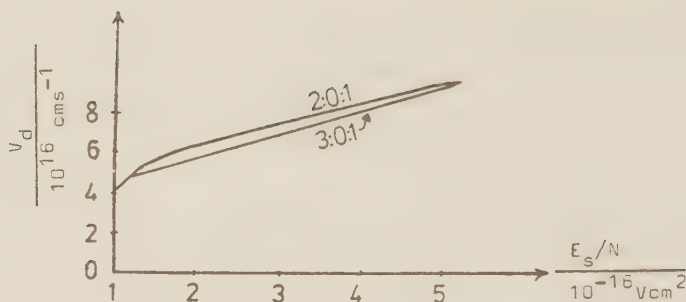


Fig.(A-5) Electron drift velocity (V_d) versus E_s/N for two different gas ratios of He,N₂ and CO₂ as numerically calculated by Judd /A25 /.

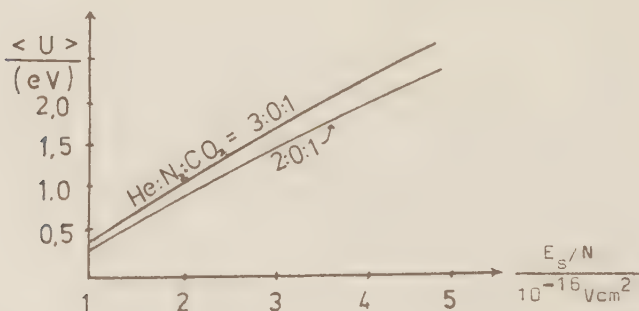


Fig.(A-6) Average electron energy ($\langle U \rangle$) as a function of the discharge parameter E_s/N for two different gas ratios /A25 /.

where, α_1 takes into account the losses due to spreading of the beam, β_1 is due to scattering and losses in the arc-protection screen, while γ_1 is the transmission factor through the window. β_1 and γ_1 are estimated to be 0.8 each for a 0.001 inch thick titanium window and gun voltages of 200-250 KV, while α_1 was in the range $0.7 < \alpha_1 < 0.85$ /A24/. Therefore,

$$\begin{aligned} J_d &= 0.5 J_p \\ &= 5.9 \text{ A cm}^{-2} \end{aligned} \quad (\text{A23})$$

The current density is related with the drift velocity of the electrons as

$$J_d = N_{eo} e V_d \quad (\text{A24})$$

where, e signifies the charge of an electron. As J_d and V_d both are known from (A23) and (A17), hence

$$N_{eo} = 6.4 \times 10^{12} \text{ cm}^{-3} \quad (\text{A25})$$

Now, the rate of electron productions (S) can be estimated with the help of (A15) as:

$$\begin{aligned} S &= \alpha N_{eo}^2 \\ &= 10^{19} \text{ cm}^{-3} \text{ s}^{-1} \end{aligned} \quad (\text{A26})$$

Numerical modelling requires not only the values of S and α but also the excitation rates. The colliding molecular species transfer energy back and forth in a variety of ways, each with its characteristic rates. These electron impact excitation rates, denoted here by X , when displayed as a function of the discharge parameter E_s/N exhibit a strong dependence on the ratio of the gases /A25/. These rates can be found in /A9, A14, A25/ and are explained in Fig.(A-2). For the above described value of U and the gas ratios these rates assume the following values, for our case:

CO_2 symmetric stretching mode,	$X_1 = 4 \times 10^{-9} \text{ cm}^3 \text{ s}^{-1}$	
CO_2 bending mode,	$X_2 = 6 \times 10^{-9}$	" (A27)
CO_2 asymmetric stretching mode,	$X_3 = 5 \times 10^{-9}$	"
N_2 vibrational mode,	$X_4 = 0$	"

A.6 De-excitation Processes

It has been mentioned in the previous section that the inelastic scattering of electrons (in case of e-beam lasers) from the gas molecules produces the population inversion. This inelastic scattering does not selectively populate the upper CO_2 laser level. It is compensated fortunately by the de-excitation times for the CO_2 molecules excited states. It can be best understood with the help of an energy level diagram for CO_2 molecules as shown in Fig.(A-6).

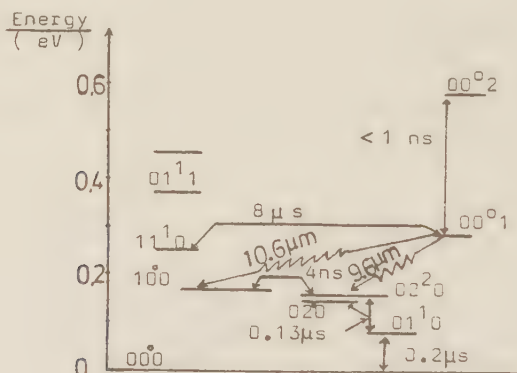


Fig.(A-6) CO_2 energy levels showing the de-excitation times for a gas composition of $\text{He}:\text{N}_2:\text{CO}_2 = 4:1:1$ /A27/. The data was extended linearly to $2.4 \times 10^5 \text{ Pa}$ on the assumption that all de-excitation phenomena are collision dominated.

It is a four level laser and the laser action takes place between the 00^0_1 and 10^0_0 levels ($10.6 \mu\text{m}$) and between 00^0_1 and 02^2_0 levels ($9.6 \mu\text{m}$) as shown in Fig.(A-6). Actually, lasing

occurs on two sets of lines centred around $10.6 \mu\text{m}$ and $9.6 \mu\text{m}$ due to the rotational levels. Patel /A26/ measured the gain as a function of J' for various transitions and found that $P(20)$ transition was the strongest one. In fact, all the transitions compete with each other and once a transition starts oscillating, it steals atoms from all of the other rotational levels by the fast relaxation among rotational levels. So, our CO_2 laser system was forced to oscillate at $P(20)$ to obtain the maximum single frequency output power.

The upper laser level (00^0_1) loses the energy by collision with the other CO_2 molecules, and may decay to the $100 \rightarrow 020$ levels by spontaneous emission. This spontaneous decay has a very long life time of about 4.2 s. Losses due to higher levels (00^0_2 etc.) are negligible as 00^0_2 level is also being pumped and it is quite possible that the equilibrium results in a net flow of population to the upper laser level (00^0_1). The advantage of using He is that the deactivation of 00^0_1 level by He is not significant but is severe by collisions with H_2 and H_2O molecules. Hence, H_2 and H_2O gases are not normally used in pulsed CO_2 lasers. So, the upper laser level is being depopulated mainly by collisions among CO_2 molecules to such as 110 level (besides the induced emission). It is essential for a laser to have a lower laser level of short life-time. The equilibrium time between $100-020$ levels is short, $\sim 4 \text{ ns}$, but they decay to 010 level with a life-time of $\sim 130 \text{ ns}$. The 010 level is the bottle-neck in the CO_2 laser with a life-time of $\sim 10 \mu\text{s}$ at a CO_2 pressure of one atmosphere. However, addition of He reduces its life-time to $0.2 \mu\text{s}$ for a gas mixture of $\text{He:N}_2:\text{CO}_2=4:1:1$ /A27/.

Excitation cross sections have been extensively studied by Nighan et al /A28, A29/. He found that the dominant loss mechanism is collision and so the loss rates vary linearly with pressure. In a TEA CO_2 laser the upper level life-time is about $10 \mu\text{s}$, the inter-rotational relaxation time is $< 1 \text{ ns}$, while the 010 level acts as a bottle neck to the system.

The above consideration leads to the conclusion that the non-selective excitation of the upper CO_2 laser level by inelastic scattering is being compensated by its long life-time and short life-time of the lower CO_2 laser level.

A.7 The Computer Model 1

In the previous sections, the excitation and the de-excitation processes have been described. The most important consideration in the design of a laser is the output energy (or power) under certain operating conditions. Now, a set of equations will be described which can be used for the explicit prediction of the output power, and population densities in various levels. This model is for a gas mixture of He, N_2 , and CO_2 . It is based on the following assumptions:

- all the rotational sublevels within each vibrational level are populated according to Boltzmann distribution,
- the translational gas temperature is obtained from the total energy in all modes and from a knowledge of the specific heat of the gas,
- the total internal losses amount to 20%, and
- losses due to attachment to neutral molecules are negligible as compared to the losses due to the recombination of electrons and ions.

The details of the model have been worked out by many authors /A12-A14/. Here, only the results will be quoted.

The following set of equations will be used to describe the time development of the energy densities stored in the vibrational modes of CO_2 , N_2 and CO together with the equation for the translational energy of the gas. Consider a schematic diagram of the set of processes which describe the $V-V$, $V-T$ transitions including the collisions processes, Fig.(A-2), E_1 , E_2 , E_3 are the symmetric, bending and asym-

metric modes of CO_2 { see Fig.(A-1) }, and E_4 and E_5 are the vibrational modes of N_2 and He. Let τ_{10} be the relaxation time for the symmetric mode, τ_{20} be the relaxation time for the bending mode and τ_3 be the V-V relaxation time, and E_1 is the energy density stored in the CO_2 symmetrical stretching mode. In this mode the energy is gained by:

- (a) photon absorption due to stimulated emission { $\nu_1 \Delta N W I_0(t)$ }
- (b) electronic excitation { $N_e(t) N_{\text{CO}_2} h \nu_1 X_1(T)$ } and

(c) due to V-V exchange among the three modes of CO_2

However, this mode loses its energy via:

- (a) V-V transitions to the bending mode
- (b) V-T transitions to the ground state

So the time development of E_1 will be /A14/:

$$\begin{aligned} \frac{dE_1}{dt} = & N_e(t) f N_{\text{CO}_2} h \nu_1 X_1 + h \nu_1 \Delta N W I_0 - \frac{E_1 - E_1^e(T)}{\tau_{10}(T)} - \frac{E_1 - E_1^e(T_2)}{\tau_{12}(T_2)} \\ & + \sum_{\lambda=3,5} \frac{h \nu_1}{h \nu_\lambda} \frac{E_\lambda - E_\lambda^e(T, T_1, T_2)}{\tau_\lambda(T, T_1, T_2)} \end{aligned} \quad (\text{A2B})$$

where,

- T = gas temperature in degrees Kelvin
- $N_e(t)$ = electron density at time t
- X_1 = effective electron vibrational excitation rate for the symmetric stretching mode
- N_{CO_2} = number density of CO_2 molecules
- ΔN = population difference between the upper and lower laser levels
- f = fraction of the undissociated CO_2 molecules
- $h \nu_1$ = energy of one symmetric stretching mode quantum
- t = time in seconds

$$E_{\lambda}^e(T, T_1, T_2) = N_{\text{CO}_2} f h \nu_{\lambda} \left[\exp \left\{ \frac{h \nu_1}{K_B T_1} + \frac{h \nu_2}{K_B T_2} + \frac{h(\nu_{\lambda} - \nu_2 - \nu_1)}{K_B T} \right\} - 1 \right]^{-1}$$

(equilibrium value of E_3 and E_5)

K_B = Boltzmann's constant

$$E_{\lambda}^e(T_j) = h \nu_1 f N_{\text{CO}_2} / \left\{ \exp(h \nu_1 / K_B T_j) - 1 \right\}$$

(equilibrium value of E_1 at temperature T_j)

$$\{\tau_{\lambda 0}(T)\}^{-1} = \text{relaxation time associated with the decay of } E_{\lambda},$$

$\lambda = 1, 2$, directly to translation

$\tau_{\lambda}(T, T_1, T_2)$ = characteristic relaxation time

$$\{\tau_{12}(T_2)\}^{-1} = \text{relaxation time associated with the transfer of energy between the CO}_2 \text{ symmetric stretching mode and the bending mode}$$

ω = stimulated emission rate on the line centre

$$= \{ F \lambda^2 4 \pi^2 h \nu_0 \tau_{\text{sp}} \Delta \nu_0 \}^{-1}$$

l' = active cavity length

l = total cavity length

F = filling factor (l'/l)

ν_0 = laser frequency

$\Delta \nu_0$ = laser transition homogeneous line width

$$= \sum_t \frac{N_t Q_t}{\pi} \left(\frac{8 K_B T}{\pi \mu_t} \right)^{1/2}$$

$$\mu_t = \frac{m_{\text{CO}_2} m_t}{(m_{\text{CO}_2} + m_t)} \quad (\text{reduced mass})$$

Q_t = collision cross-sections for species t

I_0 = laser intensity

τ_{sp}^{-1} = spontaneous transition rate

Similarly, E_2 is the energy density stored in the doubly degenerate bending mode. This mode gains energy via:

- (a) V-V transitions with the symmetric mode τ_{12} ,
- (b) V-V transitions among the three modes of $CO_2(\tau_3)$, and
- (c) electronic excitations $\{N_e(t) N_{co_2} h\nu_2 X_2(T)\}$.

It losses energy by V-T transitions to the ground state τ_{20} .

Hence the time development of E_2 , similar to E_1 , will be:

$$\begin{aligned} \frac{dE_2}{dt} = & N_e(t) f N_{co_2} h\nu_2 X_2 - \frac{E_2 - E_2^e(T)}{\tau_{20}(T)} + \frac{E_1 - E_1^e(T_2)}{\tau_{12}(T_2)} \\ & + \sum_{\lambda=3,5} \frac{h\nu_2}{h\nu_\lambda} \frac{E_\lambda - E_\lambda^e(T, T_1, T_2)}{\tau_\lambda(T, T_1, T_2)} \end{aligned} \quad (A29)$$

where T_2 = effective temperature of the CO_2 bending modes

$$= h\nu_2 / [K_B \log\{2 h\nu f N_{co_2} / E_2 + 1\}]$$

X_2 = effective electron vibrational excitation rate for the bending mode

$h\nu_2$ = energy of one bending mode quantum

E_3 represents the energy density in the CO_2 asymmetric mode, which contains the upper laser level (001). This mode gains the energy from:

- (a) the nitrogen molecule (τ_{43})
- (b) the He molecule, and from the
- (c) electronic excitations, $\{N_e(t) f N_{co_2} h\nu_3 X_3\}$.

It losses the energy by the V-V process (τ_3) and by photon emission. Therefore, we have

$$\begin{aligned} dE_3/dt = & f N_{co_2} N_e(t) h\nu_3 X_3 - \frac{E_3 - E_3^e(T, T_1, T_2)}{\tau_3(T, T_1, T_2)} + \frac{E_4 - E_4^e(T_3)}{\tau_{43}(T)} \\ & + \frac{h\nu_3}{h\nu_5} \frac{E_5 - E_5^e(T, T_3)}{\tau_{53}(T, T_3)} - h\nu_3 \Delta N W I_0 \end{aligned} \quad (A30)$$

The energy density E_4 stored in the N_2 molecule is generated by electron excitation $\{N_e(t) N_{N_2} h\nu_4 X_4\}$, and by V-V exchange among the He and N_2 molecule. It loses the energy by pumping the CO_2 asymmetric mode (τ_{43}). Therefore,

$$dE_4/dt = N_e(t) N_{N_2} h\nu_4 X_4 - \frac{E_4 - E_4^e(T_3)}{\tau_{43}(T)} + \frac{h\nu_4}{h\nu_5} \frac{E_5 - E_5^e(T, T_4)}{\tau_{54}(T, T_4)} \quad (A31)$$

E_5 refers to the energy density stored in the CO vibrational mode which is produced when a fraction of the CO_2 molecules dissociates into CO. Electrons will excite the CO molecule into the first excited state, $v=1$. The resulting excited CO molecule will then be available to transfer energy in V-V collisions to both N_2 and CO_2 . Hence

$$\frac{dE_5}{dt} = N_e(t)(1-f)N_{CO_2} h\nu_5 X_5 - \sum_{\lambda=3,4} \frac{E_5 - E_5^e(T, T_\lambda)}{\tau_{5\lambda}(T, T_\lambda)} - \frac{E_5 - E_5^e(T, T_1, T_2)}{\tau_5(T, T_1, T_2)} \quad (A32)$$

These equations represent the rate of flow of energy in and out of the vibrational modes. Vibrational mode energy could be transferred to the kinetic energy of the molecules and may increase the ambient temperature, as both are related by

$$(1/2)Mv^2 = (3/2)K_B T$$

where an energy of $(1/2)K_B T$ is associated with each degree of translational freedom. It amounts to total of $(3/2)K_B T$ as the motion is in three dimensions. N_2 and CO_2 are both linear molecules and may have two degrees of freedom. So, the total gas kinetic energy per unit volume is

$$E_t = (5/2)K_B T(N_{N_2} + N_{CO_2}) + (3/2)K_B T(N_{He}) \text{-----} (A34)$$

Its time rate of change may be written as

$$\begin{aligned} \frac{dE_t}{dt} = & \sum_{\lambda=1,2} \frac{E_\lambda - E_\lambda^e(T)}{\tau_\lambda(T)} + \sum_{\lambda=3,5} \left(1 - \frac{h\nu_2}{h\nu_\lambda} - \frac{h\nu_1}{h\nu_\lambda}\right) \frac{E_\lambda - E_\lambda^e(T, T_1, T_2)}{\tau_\lambda(T, T_1, T_2)} \\ & + \sum_{\lambda=3,4} \left(1 - \frac{h\nu_\lambda}{h\nu_5}\right) \frac{E_5 - E_5^e(T, T_3)}{\tau_{5\lambda}(T, T_3)} \end{aligned} \quad (A35)$$

To complete the kinetic model, an intensity equation is required which can be developed by the following considerations: If $N_{001}P(J')\{N_{100}P(J'+1)\}$ are the number of CO_2 molecules per unit volume in the rotational level $J'\{J'+1\}$ of the lowest symmetric {asymmetric mode} and $\varphi(\nu)$ is the number of photons per unit volume, then the net stimulated increase in radiation energy density per unit time is:

$$\begin{aligned} & h\nu \{ N_{001}P(J')B_{31} - N_{100}P(J'+1)B_{13} \} \varphi(\nu) \\ & = h\nu \left(\frac{c^3}{8\pi h\nu^3 \tau_{sp}} \right) \Delta N \varphi(\nu) \end{aligned} \quad (A36)$$

where, the first term represents the stimulated emission process and the second term represents absorption per unit volume per unit time, $B_{nm} (= \frac{c^3}{8\pi h\nu^3 \tau_{sp}})$ is the Einstein coefficient,

$P(J')$ is the Boltzmann distribution factor.

By defining the radiation flux as $c\varphi(\nu) = I_\nu \delta(\nu - \nu_0)$, (A36) can be written as

$$\frac{\Delta N c^2}{8\pi \nu^2 \tau_{sp}} \{ I_\nu \delta(\nu - \nu_0) \} g(\nu) d\nu \quad (A37)$$

where, $\delta(\nu - \nu_0)$ is the dirac delta function to characterize the narrow laser line width, and $g(\nu)d\nu$ is the probability of a transition between the energy levels resulting in an emission or an absorption of a photon whose energy is in the range $h\nu$ and $h(\nu + d\nu)$. An integration of (A37) over the frequency yields:

$$\frac{\Delta N c^2 I_{\nu_0} g(\nu_0)}{8\pi \nu_0^2 \tau_{sp}} \quad (A38)$$

So, the rate of increase of radiation flux is

$$\begin{aligned} \frac{dI_{\nu_0}}{dt} &= \frac{\Delta N c^3 I_{\nu_0} g(\nu_0)}{8\pi \nu_0^2 \tau_{sp}} \\ &= \Delta N c h \nu_0 \omega I_{\nu_0} \end{aligned} \quad (A39)$$

where, $g(v_0)$ is replaced by $\frac{2}{\pi \Delta \nu}$ (for Lorentz frequency distribution). On the other hand, all the losses, for instance, transmission, absorption, and scattering by the mirrors, can be accounted for by a single parameter τ_c . τ_c is generally known as the life-time of a photon within the laser cavity, and is

$$\tau_c = - \frac{(2L'/c)}{\ln(1/\delta)} = - (2L'/c) \ln R$$

where L' is the length of the laser cavity, and $R=1-\delta$.

In the absence of the spontaneous emission, the net increase in the radiation intensity in the cavity will be:

$$\frac{dI_{v_0}}{dt} = - \frac{I_{v_0}}{\tau_c} + c h v_0 \Delta N I_{v_0} \quad (A40)$$

$$\text{where, } W = \left(\frac{\lambda_0^2}{4 \pi^2 h v_0 \tau_{sp}} \right)$$

The spontaneous emission term is given by Gilbert et al /A28/

$$\frac{c h N_{001} P(J) v_0}{\tau_{sp}} G \quad (A41)$$

where G is the fraction of spontaneously emitted photons into the small angular aperture of the resonator within the spectral width $d\nu$ of a single axial mode. If Ω is the solid angle then the fraction of radiation retained within the cavity will be $\Omega/4\pi$ within a spectral width of $g(v_0)d\nu (= \frac{2}{\pi} \frac{d}{\Delta \nu})$.

$$\text{Hence, } G = \frac{\Omega}{2\pi^2} \frac{d\nu}{\Delta \nu}$$

where, $\Delta \nu \approx 4$ GHz at STP and $d\nu \approx (2\pi \tau_c)^{-1}$. Addition of the spontaneous emission in (A40) gives

$$\frac{dI_0}{dt} = - \frac{I_0}{\tau_c} + c h v_0 \Delta N I_0 + c h v_0 N_{001} P(J) U \quad (A43)$$

$$\text{where, } U = \frac{\Omega d\nu}{2\pi^2 \tau_{sp} \Delta \nu}, \quad I_{v_0} = I_0$$

(A43) gives the temporal evolution of the cavity field intensity. The knowledge of the intensity helps to evaluate the output power of the laser by the following simple considerations:

Consider a cavity of length L' filled with an active medium. Let α_g be the gain per unit length and the cavity power flux travelling towards the output mirror is y .

If δ is the geometric output coupling factor, then $(1-\delta)y$ will be reflected back by the output coupling mirror. Hence, for steady state (neglecting all other losses):

$$e^{\alpha_g L'} (1-\delta)y = y$$

$$\text{or } \alpha_g = - (1/2L') \ln(1-\delta) \quad (\text{A44})$$

From Fig.(A-8), it is self explanatory that the output intensity is:

$$I_o = y\delta \quad (\text{A45})$$

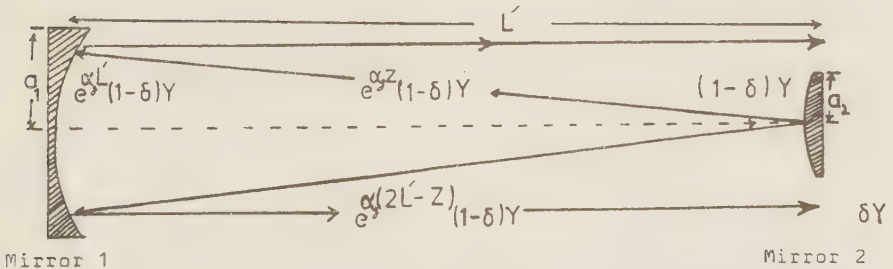


Fig.(A-8) Variation of the cavity field intensity in a positive branch confocal unstable resonator is schematically shown here. If y is the power flux travelling towards the output mirror and δ is the output coupling factor then the intensity coming out of the resonator will be δy of which a proportion $(1-\delta)y$ will be reflected back and get amplified, where α_g is the small signal gain coefficient.

If A is the cross-sectional area of the output ring, then

$$P_o = Ay\delta \quad (A46)$$

The kinetic equations yield the information about the intensity I_o , but not y . However, a relationship between them can be found by defining

$$\begin{aligned} I_o &= \frac{1}{L} \int_0^L I_z dz \\ &= y\delta / \alpha_1 L \end{aligned}$$

where, I_z is the laser intensity at a certain point, say z within the cavity and will be the sum of the amplified intensity being reflected from both the mirrors. By substituting the value α_1 from (A44), y can be written as:

$$y = \frac{1}{2} \frac{\ln(1-\delta)}{\delta} I_o \quad (A47)$$

The knowledge of y was essential to determine the output power, which in our case now can be written using (A46) and (A47):

$$P_o = - \frac{A}{2} \ln(1-\delta) I_o \quad (A49)$$

A.8 Comparison Of Computer Results With Experiment

Equations (A28-A35) were solved along with the intensity equation (A43) by applying some numerical techniques, like Runge-Kutta method. A typical computer (Cyber 175) result for the input parameters, described in the previous sections, is shown in Fig. { (A-9) - (A-17) }. The horizontal scale is in microseconds and $t_e = 0$ corresponds to the time when e-beam gun starts emitting electrons as shown in Fig.(3,8). In these figures, the relevant population densities of upper laser level (Fig.A-10) and the population inversion density


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.26091-09 .55071+00 .14651+04 .63391+01 .74241-02 .12071+17 .14241+17 .21701+16 325.1 .311331+02 .912271+07
.26101-09 .55081+00 .14701+04 .63431+01 .74301-02 .12101+17 .14271+17 .21671+16 325.2 .311331+02 .912311+07
.26111-09 .55091+00 .14751+04 .63471+01 .74351-02 .12131+17 .14301+17 .21641+16 325.3 .311331+02 .912351+07
.26121-09 .55101+00 .14801+04 .63511+01 .74401-02 .12161+17 .14331+17 .21611+16 325.4 .311331+02 .912391+07
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.26151-09 .55131+00 .14951+04 .63631+01 .74551-02 .12251+17 .14421+17 .21521+16 325.7 .311331+02 .912511+07
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.26321-09 .55301+00 .15801+04 .64311+01 .75401-02 .12761+17 .14931+17 .21011+16 327.4 .311331+02 .913191+07
.26331-09 .55311+00 .15851+04 .64351+01 .75451-02 .12791+17 .14961+17 .20981+16 327.5 .311331+02 .913231+07
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.27001-09 .55981+00 .19201+04 .67031+01 .78801-02 .14801+17 .16971+17 .18971+16 334.2 .311331+02 .915911+07

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Fig.(A-9) A typical computer model 1 result with the input parameters as described in sections (A.4) and (A.5). Only the initial and final results are shown as the interesting parameters have been plotted in the next figures.

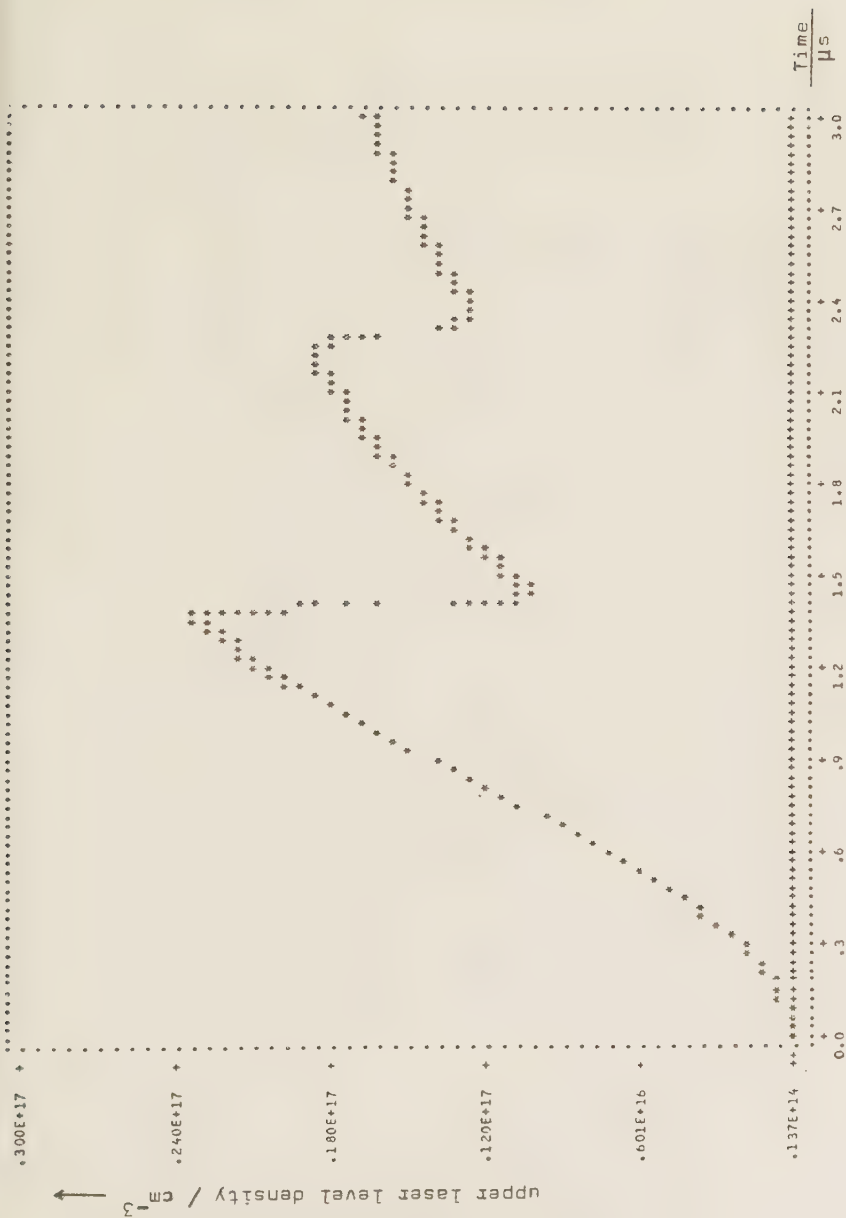


Fig.(A-10) Time development of the upper laser level (00^0_1) population density.

(Fig. A-11) are plotted as a function of time. In addition to it, the laser output power, the temperature variation in the cavity, and the energy density stored in the several modes are shown in Fig. (A-12, A-13, A-14, A-15). Several important informations can be obtained from these figures. Fig. (A-13) shows that the ambient temperature in the discharge chamber does not increase to a large extent. Hence, the thermal effects are negligible and there is no danger for the output deterioration. Next, consider Fig. (A-16) in which the population inversion density and the laser output power are plotted on the same time scale for the sake of comparison. The laser output occurs twice. The first one is the main laser pulse and the second one is a delayed laser pulse. The mechanism by which the laser output is formed may be described as follows. The laser action is initially started by the spontaneous noise and later lasing occurs when the pump excitation exceeds the threshold. At the initial stages, no lasing occurs so the population density in the upper laser level (Fig. A-10) and hence, the population inversion density (Fig. A-11) both increases rapidly. It is due to the exponential nature of the amplification which is obvious from the gain curve (Fig. A-17). The time difference between the main laser pulse and the initiation of the excitation pulse depends on the steepness of the leading edge of the excitation pulse, rate of rise of the gain and the speed of the relaxation processes. At the final stage, lasing occurs when the population inversion is achieved and the supply of the excited molecules begins to decrease due to rapid conversion to photons through stimulated emission.

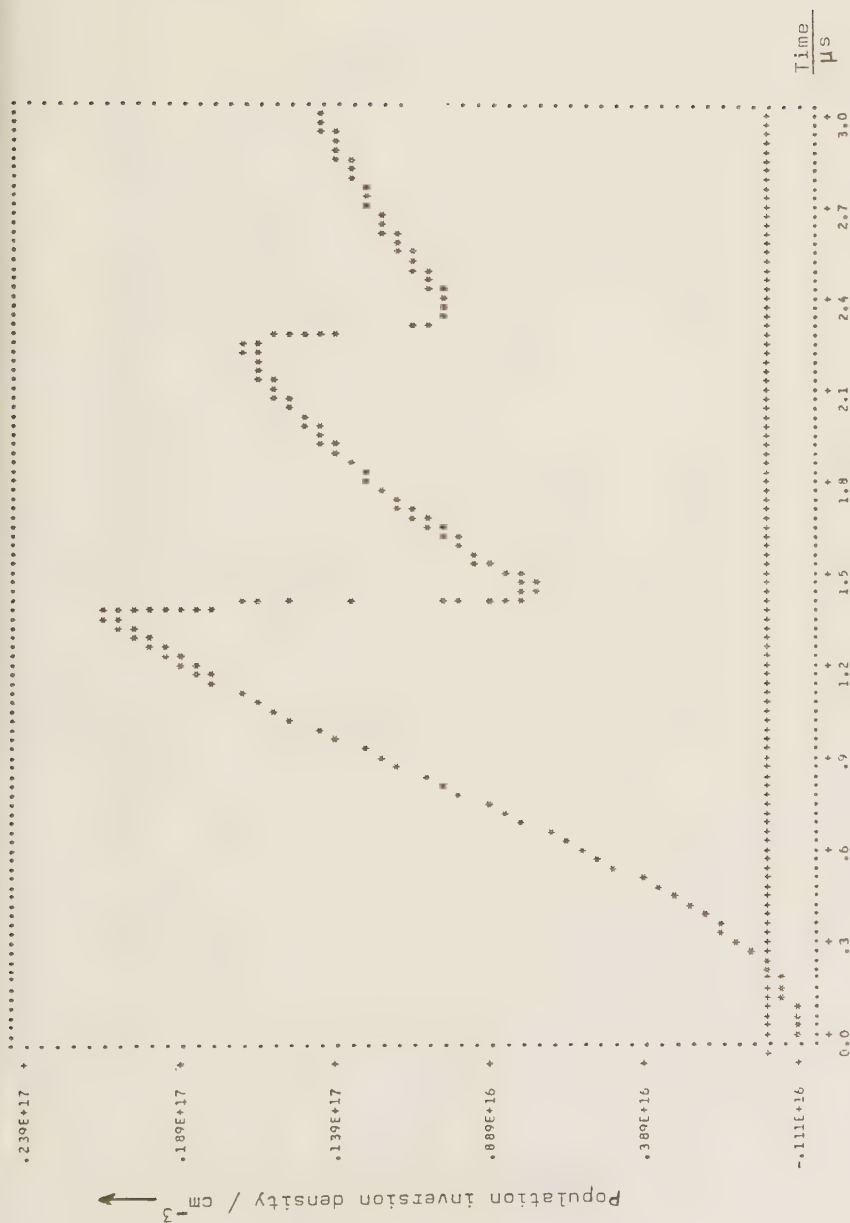


Fig.(A-11) Time development of the population inversion density.

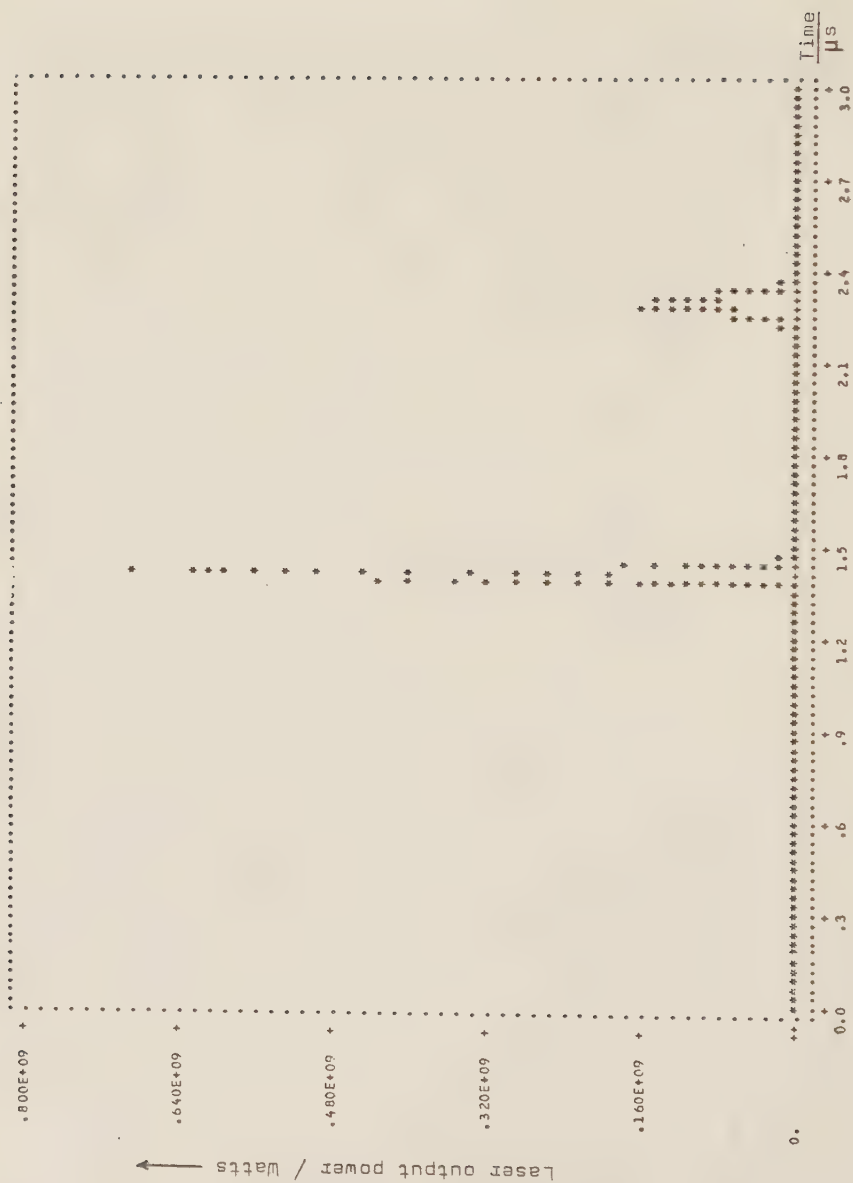


Fig.(A-12) Laser output power as a function of the time.

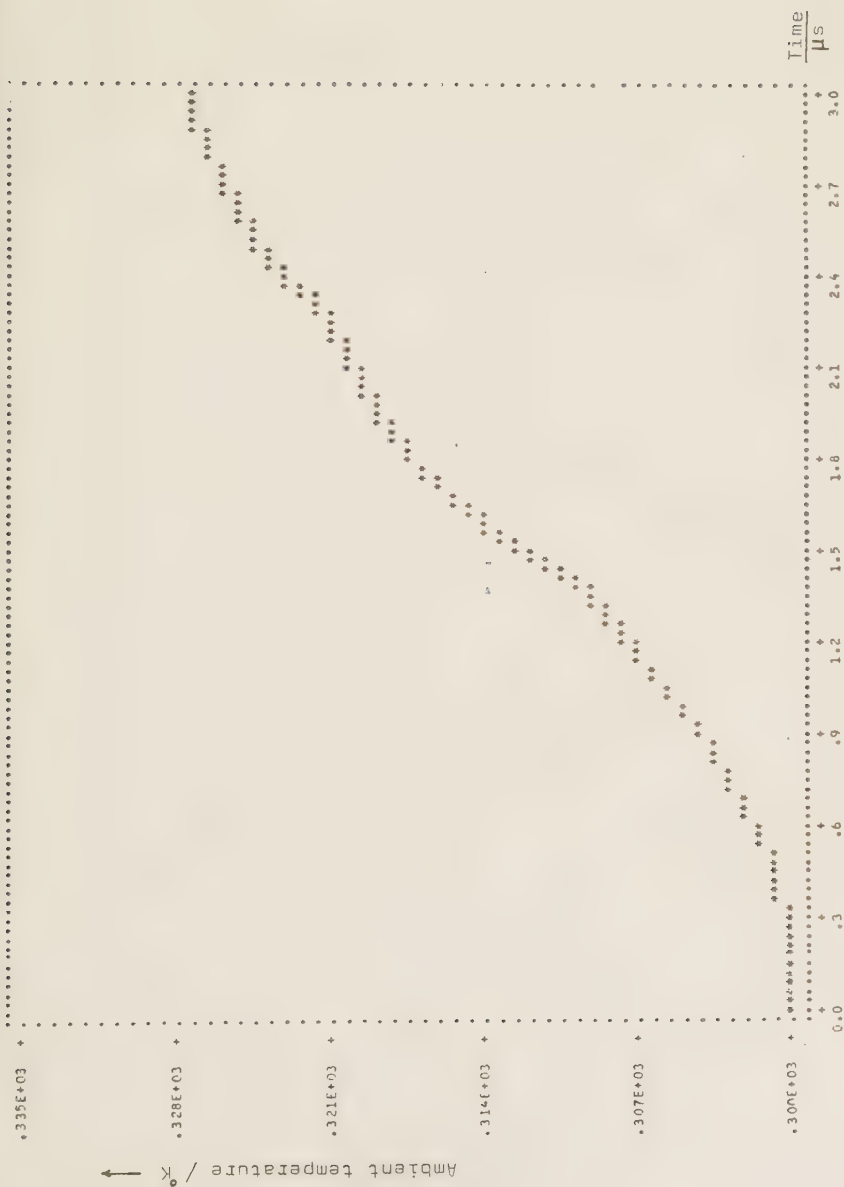


Fig. (A-13) Variation in the ambient temperature of the gas mixture with time.

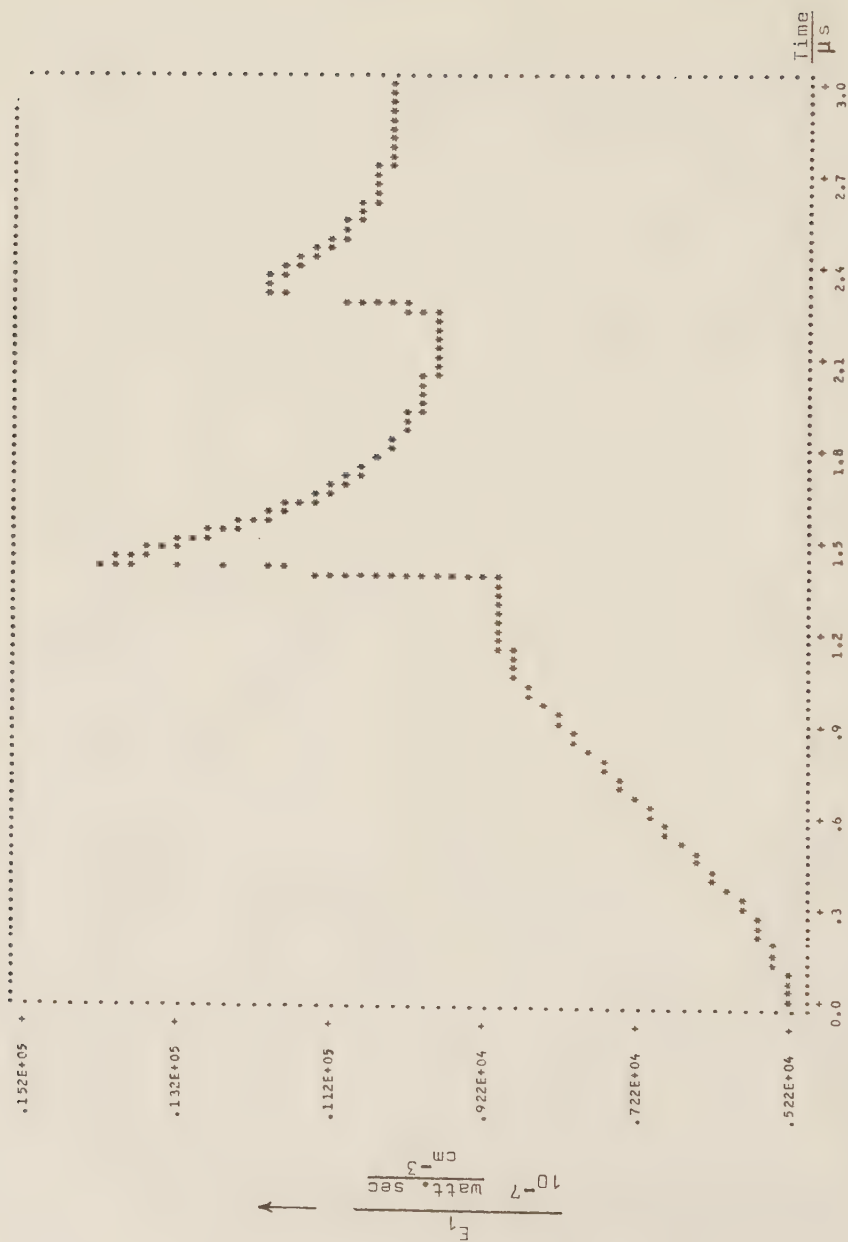


Fig.(A-14) Energy density (E_1) variation in the lower laser level with time.

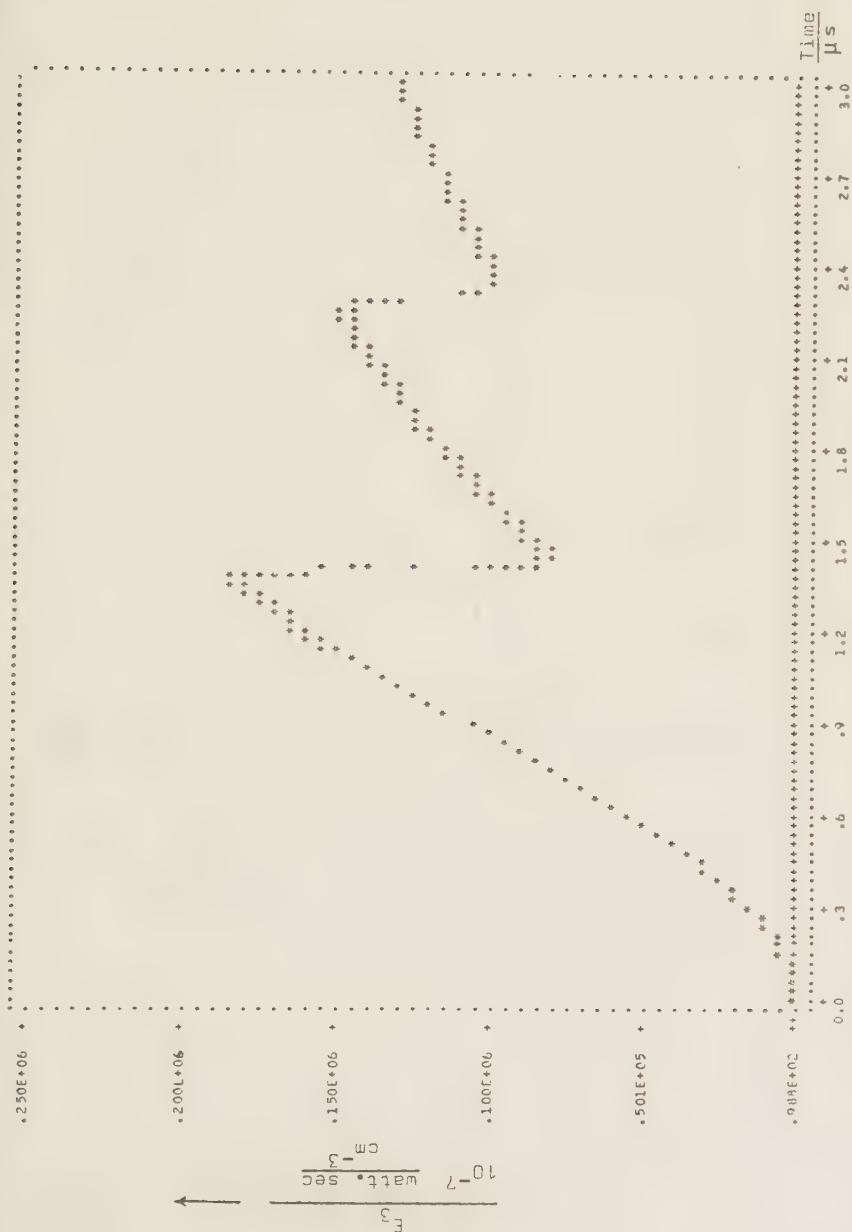


Fig. (A-15), Energy density (E_3) variation in the upper laser level with time.

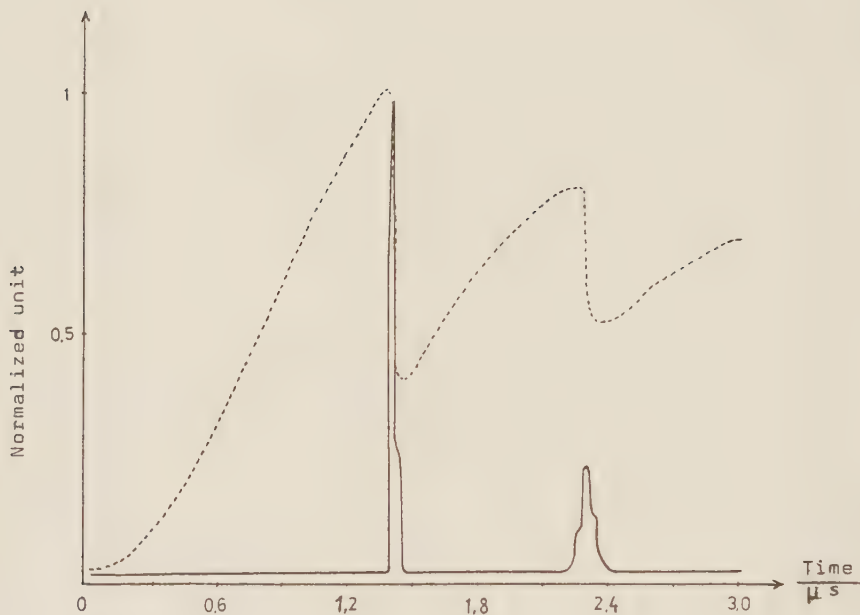


Fig.(A-16) Time dependence of the population inversion density (dotted curve) and the laser output power (solid curve) on the same scale as predicted by the computer model 1.

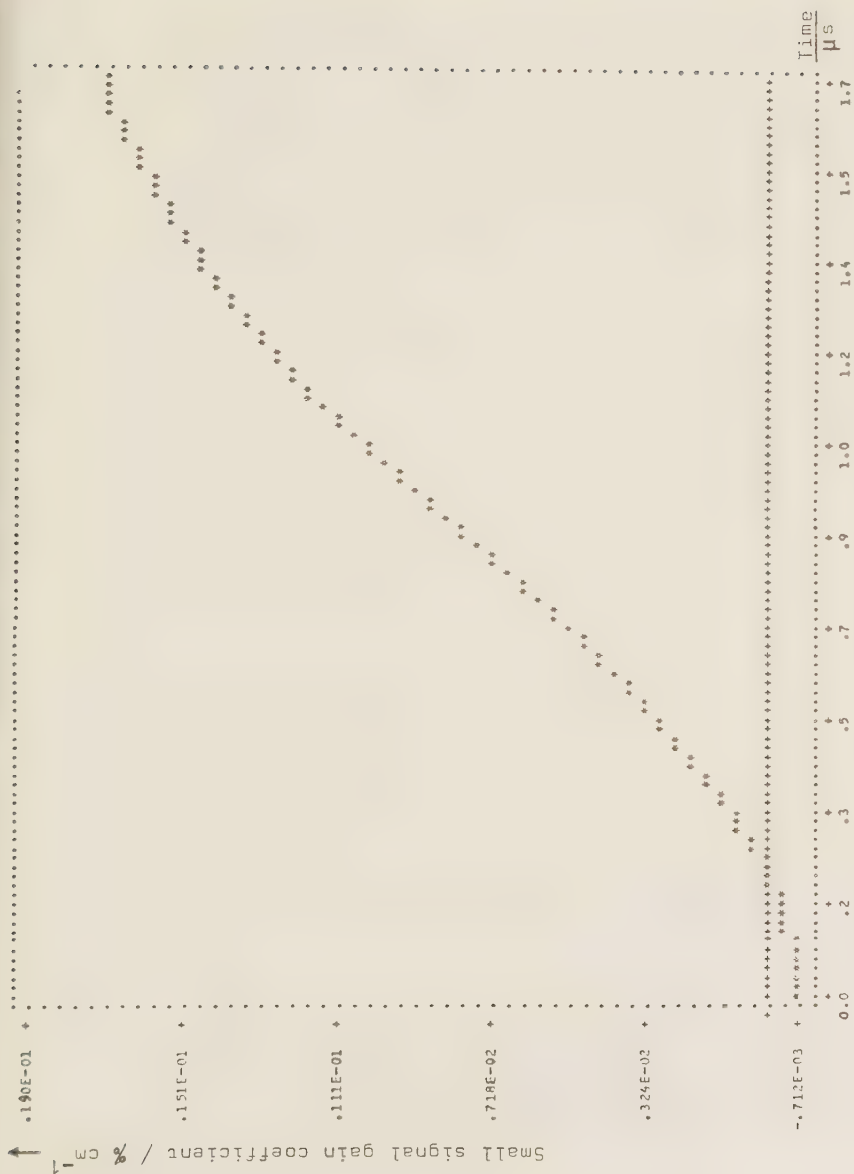


Fig.(A-17) Small signal gain coefficient as a function of time.

The population of the lower laser level will increase and even may become equal to the population density of the upper laser level due to a sharp drop in the population density of the upper laser level. Consequently, it results in a sharp drop in the population inversion density, as obvious in the figures (A-10, A-11). Although the gain has not decreased but it is not sufficient to make up for the dissipative losses and the laser pulse collapses shortly. The main laser pulse consumes only a fraction of the total energy stored in the populations, as obvious from Fig.(A-14, A-15). As the gain has not reduced and the lower laser level has a fast relaxation, population inversion again builds up and again lasing occurs which gives a delayed laser pulse. As the lower laser level was still excited, so the power of the delayed pulse was less than that of the main laser pulse. All these computer predictions were found to be in good agreement with the experimental observations. N_2 gas was not, therefore, deliberately added to reduce the energy of the delayed pulse. Since delayed pulses were not required for a LPI experiment but they may be harmful as they can damage the optical components and may complicate the situation.

The computer model results will be checked by a quantitative comparison with our experimental results. The computer model predicts an output power of about 700 MW { Fig.(A-12) } and an output energy of 5.286 J/l { Fig.(A-9) } . The active volume of the laser was 7.85 litres. Hence, it corresponds to a total output energy of about 42 J. It is in good agreement with the measured values of total output energy of about 45 J and an output power of $\sim$ 500 MW. The small discrepancies can be easily accounted for by the uncertainties in the absolute calibration of the detectors.

The above considerations show that a good correlation exists between the experimental results and the computer model predictions. On this basis, the computer model can be used to see the effect of the various parameters on the laser output energy (or power). They are:

- If the geometric output coupling parameter (δ) of the resonator (section A.4) is varied, then the output power of the laser also changes. Such an effect has been shown in Fig.(A-18), where the computer model results for different values of δ are plotted, while all other parameters are being kept constant. It shows an optimum value of δ as 0.75, which confirms that our laser resonator design was optimum in this regard.

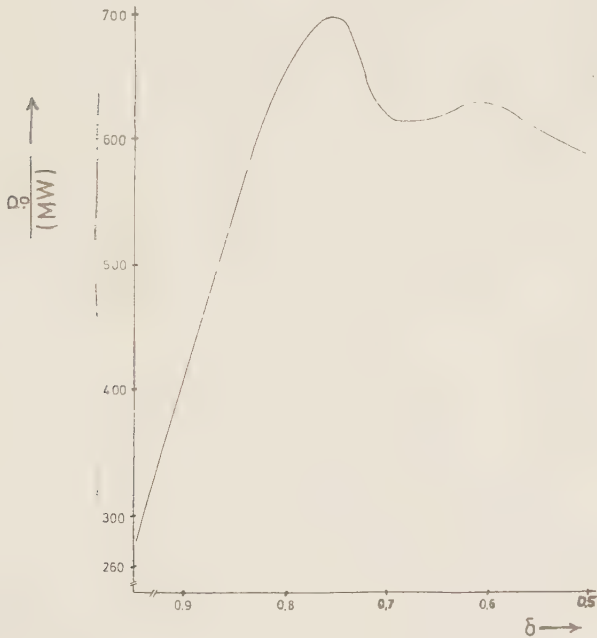


Fig.(A-18) Variation of the laser output power, P_o , as a function of the geometric output coupling parameter, δ , as predicted by the computer model 1.

- In a CO_2 laser system, the output energy depends strongly on the ratio of the gases, as shown in Fig.(A-19). In this figure, both the experimental results and the computer results are shown on the same scale. The computer results are partially in good agreement with the experimental results for CO_2 in the range of 10-30%. For CO_2 ratio of 10% and 45% in the gas mixture computer model predicts no output. On the other hand, experimentally some laser output was observed for 10 and 45% of CO_2 but it was highly unreliable. The difference between the experimental results and the computer model predictions could be explained on the basis of the variation in the values of the excitation (kinetic) rates. They depend on the gas ratio and on the discharge conditions. The discharge conditions deteriorate for very low (10%) and for high values (50%) of He, and it impedes the estimation of the kinetic rates. As the experimental results are in good agreement with the computer results within a certain range, it can be concluded that the optimum value of CO_2 was 30%.

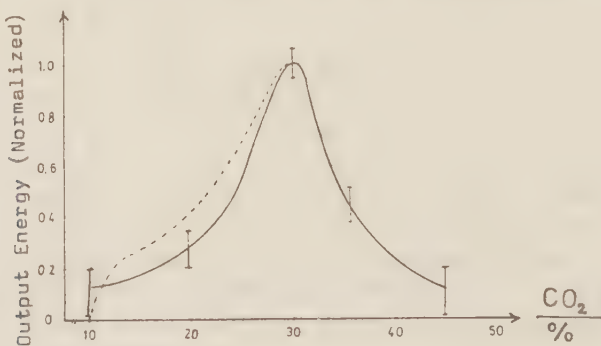


Fig.(A-19) Effect of the CO_2 ratio in the gas mixture on the laser output energy. Dotted curve is the computer model prediction while the solid curve is the experimental results.

- In our resonator, there are certain internal losses, for instance, imperfect reflection of the mirrors (say, 95% each) radiation coming out of the hole in the rear mirror (say, 5%), and reflection from the coupling window (4% from each surface). It amounts to a total internal loss of about 20%. The computer model predicts (Fig.A-20) a decrease in the output power as the internal losses increase beyond 15%. If the internal losses are 20%, then the laser output power by about 7% reduces. However, a further 5% increase in the internal losses (i.e. 25% internal losses) reduces the laser output power by an amount of about 17% which is quite significant. Hence, it shows that not much power was lost in our case.

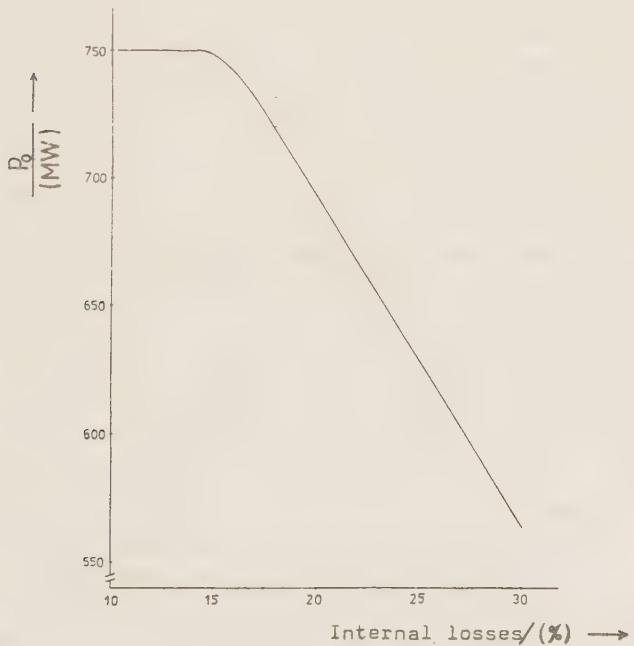


Fig.(A-20) Computer model prediction for the effect of the internal losses on the output power.

A.9 The Computer Model 2

As described in the previous section, the computer model 1 predictions help to optimize the laser system. In addition to it, this computer model furnishes the information about the time development of the upper and lower laser level which are essential for the description of IML by another computer model number 2.

For IML, a short laser pulse of a few nanoseconds is fed into the slave oscillator which after many passes results in a large amplified pulse. In the previous description of computer model 1, it was assumed that the rotational relaxation is instantaneous which implies that the population distribution in all the rotational sublevels within each vibrational level is in accordance with the Boltzmann law. This is true in the case of long laser pulses like one which we have without IML. On the other hand, this assumption is not valid in case of short laser pulses which results due to IML at atmospheric pressure.

It is important to know when the rotational effects become significant. It can be estimated as: Cheo et al /A29/ have measured the rotational relaxation constant k_{rot} of the CO_2 00^0_1 upper laser level as $(1.1 \pm 0.3) \times 10^7 \text{ s}^{-1} \text{ torr}^{-1}$. The presence of the He affect k_{rot} and it reduces this constant to $(0.7 \pm 0.1) \times 10^7 \text{ s}^{-1} \text{ torr}^{-1}$ /A30/, which corresponds to a rotational relaxation time of 94 ps in our case for a gas pressure of $2 \times 10^5 \text{ Pa}$. It was found /A31, A32/ that even for laser pulses of up to 10^2 times, the rotational coupling times, the coupling between the rotational energy levels, has a significant impact on laser dynamics. Hence, this limit is about 10 ns in our case. It justifies our assumption of neglecting the rotational effects in computer model 1 where the FWHM of the laser pulse was 70 ns but is no more negligible for laser pulses of about 2 ns durations as we have in the case of IML.

The saturation characteristics of high pressure molecular laser amplifiers are significantly affected by rotational energy transfer. It results in a substantial increase in the

energy gain experienced by a few nanosecond duration pulse in a high pressure CO_2 laser amplifier. Therefore, a new computer program (modified Pulsamp/A15/) will be employed for the description of IML in our case. This computer model 2 is based on the following assumptions:

- the population inversion between the two vibrational levels of the amplifying process is not increased by inter level relaxation or repumping since these processes are negligible in the CO_2 , He, N_2 kinetics at a pressure of 10^5 Pa and on a 5 ns scale /A31/,
- the laser medium consisted of a prepumped gas mixture and the population densities in the upper and lower laser level for different times at t_e can be used from computer model 1, where $t_e=0$ is the instant when the e-beam gun fires,
- input and output pulse are both single mode (P-2L),
- all the rotational sublevels are initially in thermal equilibrium with each vibrational level, and
- the input pulses are of a few nanosecond durations and the gas pressure is high (equal or above atmospheric pressure) hence, the effect of the coherent propagation of pulse is neglected /A15/. Therefore, the model is not valid for the input pulse of durations less than the dipole dephasing time, which is about 60 ps at atmospheric pressure.

The computer model 2 essentially solves a set of rate equations of Schappert /A31/. If a pulse of frequency ν_0 and intensity $I_{\nu_0}(=I_0)$ propagates through a prepumped medium, then its amplification due to the single $J'_0 \rightarrow J'_0+1$ transition can be described by

$$\frac{\partial I_0}{\partial t} + \frac{\partial I_0}{\partial x} = h\nu_0 \delta^{J'_0} \omega I_0 \quad (\text{A50})$$

and the time development of the population inversion density is:

$$\frac{\partial \delta^{J'_0}}{\partial t} = \delta^{J'_0} \omega I_0 \left(1 + \frac{g^{J'_0}}{g^{J'_0+1}}\right) - \frac{\delta^{J'_0} - P(J'_0)}{\tau_R} \quad (\text{A51})$$

$$\frac{\partial \delta^{J'}}{\partial t} = \frac{-\delta^{J'} - P(J') \Delta}{\tau_R}, \text{ all } J' \neq J'_0 \quad (\text{A52})$$

where, $n_j^{u(1)}$ is the population density of the j th rotational sublevel of the upper (lower) laser level, τ_R is the rotational time, and

$$\Delta = \sum_j n_j^u - \sum_j n_j^l = N^u - N^l$$

The input pulse will stimulate the transition between the J' (upper) and $J'+1$ or $J'-1$ (lower) sublevels whenever the energy difference between the two levels is $h\nu_0$. If this is the case for J'_0 (upper) and J'_0+1 (lower) levels, then in equilibrium, the population difference will be

$$\delta^{J'} = P(J') N^u - P(J'+1) N^l \left(\frac{\mathfrak{g}_{J'_0}^{J'}}{\mathfrak{g}_{J'_0+1}^{J'}} \right) \quad (\text{A53})$$

where, $\mathfrak{g}_{J'_0}^{J'} (= 2J'_0+1)$ is the degeneracy of the J'_0 th rotational sublevel, and, if the rotational levels are in thermal equilibrium at a temperature T then each $N_{J'}$ is a fraction $P(J')$ of the total N , and

$$P(J') = \mathfrak{g}_{J'} \exp\left(-\frac{\epsilon(J')}{k_B T}\right) \left[\sum_{J'} \mathfrak{g}_{J'} \exp\left(-\frac{\epsilon(J')}{k_B T}\right) \right]^{-1}$$

with $\epsilon(J')$ as the energy of the J' th sublevel.

The above mentioned equations were solved by using a Simpson integration formula with the help of a Cyber 175 computer. This computer model 2 predicts the behaviour of a laser pulse as it propagates in an amplifier or in a slave oscillator. Fig.(A-21) shows the e-beam laser output intensity when a laser pulse of intensity 10^{15} W m^{-2} and of FWHM of 2 ns was injected in its resonator (Fig.A-8) after 500 ns of the initiation of the e-beam gun firing. As the laser pulse is injected at $t_e=500 \text{ ns}$ along the axis of the resonator, there will not be any output at this instant, but after a round trip time the amplified pulse will be partially transmitted. The output coupling factor is very high i.e. 75% and the gain of the medium at the initial stage is not sufficient to compensate it; so first the output intensity will reduce. The gain of the medium as shown in Fig.(A-17) increases

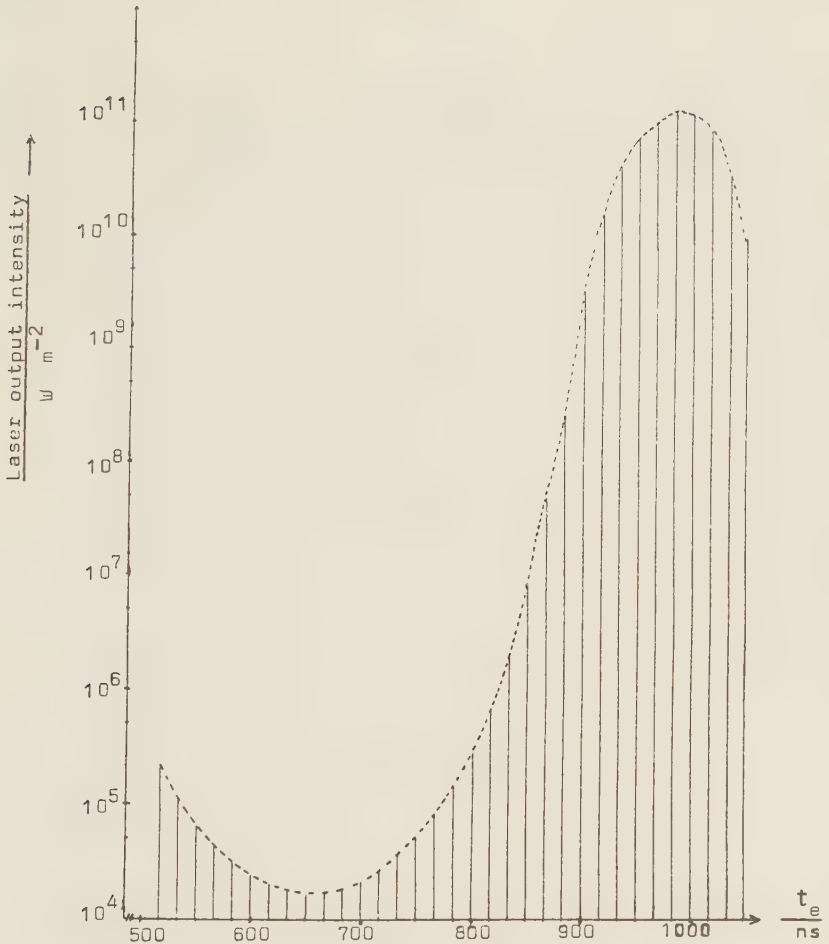


Fig.(A-21) Laser output intensity as a function of time when a laser pulse of intensity 10^5 W m^{-2} and of FWHM of 2 ns was injected at $t_e = 500 \text{ ns}$ in the unstable resonator with a geometric output coupling of 75%, where t_e is the time when the e-beam gun starts emitting electrons. One fourth of the output intensity is reflected back by the front mirror in the resonator and is being amplified again. The round trip time of the resonator is 16.7 ns. Each vertical line indicates a narrow laser pulse.

exponentially with the time and consequently the laser output of the slave oscillator increases and reaches a maximum (Fig.A-21). At this stage the population density of the lower laser level starts increasing and thus the output intensity decreases. This prediction of computer model 2 is in good agreement with the experimental observations as shown in Fig.(3.7c). In this picture, all the laser output pulses can not be recorded as the range of the detection system was about an order of magnitude. Only the laser output about the time of maximum power was recorded as it was of much interest for LPI experiment.

Fig.(A-21) shows that a laser pulse can be amplified to six orders of magnitude over a time of about 500 ns in an amplifier resulting in a much higher output intensity.

The above mentioned consideration is for a single line operation, however, it can be extended easily to take into account multiline operation, too.

List of Symbols

Here, a list of symbols has been given which are used in the text and in the Appendix. In brackets page numbers are given where these are introduced and defined.

$\vec{A}$	vector potential (11)
$\vec{A}_0(-)$	vector potential of the incident (back scattered) em- radiation (13)
A	area of cross section of the laser beam in plasma (29)
a_i	radius of the i th mirror (144)
B	initial noise level (45)
δB^h	high frequency magnetic field (8)
B_{nm}	Einstein coefficient (161)
c	velocity of light (8)
c_s	phase velocity of the ion acoustic wave (8)
d	anode cathode spacing (50)
E_1, E_2, E_3	energy density stored in the symmetric, bending and asymmetric modes of CO_2 (156)
E_4	energy density stored in the vibrational mode of N_2 (159)
E_5	energy density stored in the CO vibrational mode (160)
E_t	total gas kinetic energy per unit volume (160)
E_o	electric field of the pump wave (8)
E_{th}	threshold electric field of the ion acoustic wave (24)
E_s	sustainer voltage (54)
$f(v_i)$	ion velocity distribution function (33)
$(f^1)f$	(low frequency) pondermotive force (7)
g	ratio between the damping rate and bounce frequency (37)
g_i	curvature parameter (144)
G	fraction of the spontaneously emitted photons (162)
H	collision parameter (41)
$i=1$	represents the rear mirror (144)
$i=2$	represents the output coupling mirror (144)
I_o	laser intensity (21)
$I(\psi)$	far field intensity distribution (147)
δI	laser intensity perturbation (26)
$j=e$	stands for electron (9)
$j=i$	stands for ion (9)

J	particle current density (11)
δJ^h	high frequency electron current in the plasma (8)
J'	rotational quantum number (143)
J_0	zeroth order Bessel function (148)
J_d	current density entering in the laser discharge chamber (151)
J^p	total current density flowing from a planar cathode (151)
$\vec{k}$	wave number of an ion-acoustic wave in the plasma (14)
$\vec{k}_0(-)$	wave number of the incident (scattered) radiation (14)
$k_{\max}$	maximum value of k (16)
k_{de}	inverse of the electron Debye length (20)
k_1, k_2	wave numbers of the daughter ion acoustic wave (24)
k_{th}	thermal conductivity coefficient (29)
K_B	Boltzmann constant
$K_B T_e$	electron temperature in eV (28)
$K_B T_0$	initial ion temperature (44)
$K_B T_i$	ion temperature in eV (44)
l	degeneracy of the vibrational levels (142)
$\ln \Lambda$	Coulomb logarithm (28)
l'	active cavity length (158)
L'	distance between the mirrors in a laser resonator (145)
L	length of the plasma slab (18)
M	geometrical magnification factor (146)
M_1	rear mirror of the resonator (52)
M_2	front mirror of the resonator (52)
m_j	mass of the j th type particle (9)
N_j	density of the j th type particle (11)
n_0 or n	plasma density (9/28)
δn	density perturbation (7)
n_{cr}	critical density of the plasma (18)
ΔN_t	excess trapped particle density (40)
ΔN	population inversion density between the upper and lower laser level (157)
N	gas density (150)
N_{if}	Fresnel number (144)
N_{eq}	equivalent Fresnel number (146)

P_o	laser output power (56)
P_o/A	laser intensity (30)
P_l	power lost by the ion acoustic wave (35)
P_B/P_o	relative back scattering power (79)
P_m	maximum power of the ion acoustic wave according to Manley-Rowe relation (35)
$P(J')$	Boltzmann distribution factor (161)
Q	local rate of heat flux absorption (28)
q_j	charge of the j th type particle (11)
q	heat flux (29)
R	amount of back scattering ($=P_B/P_o$) (20)
R_i	radius of curvature of the i th mirror
S	rate of production of electrons and ions (149)
T_j	temperature of the j th type particle (9)
T	gas temperature in degrees Kelvin (157)
$\Delta t'$	time difference between the injected signal entering the slave oscillator and its emergence (62)
Δt	time difference between the incident laser and the plasma maximum compression (83)
t_e	time at which e-beam gun starts emitting electrons (164)
$\langle u \rangle$	average electron energy (151)
v_d	electron drift velocity (151)
v_o	electron quiver velocity (8)
v_j	thermal velocity of the j th type particle (8)
v_{im}	maximum velocity gained by the ions (33)
w	stimulated emission rate (158)
w_1	energy gained by the ions (32)
w_2	energy gained by the ion acoustic wave (32)
Δw	energy lost by the ion acoustic wave (33)
x	excitation rates (142)
Y	cavity power flux travelling towards the output mirror (163)
Z	charge number (8)

α_h	absorption coefficient (28)
α_1	losses due to spreading of the beam (153)
α	recombination coefficient (150)
α_g	gain per unit length (163)
β_0	ratio between the plasma density to the critical density (21)
β	attachment coefficient (150)
β_1	losses due to scattering (153)
γ_1	transmission factor (153)
$\gamma_{SBS/SRS}$	growth rate of SBS / SRS (15/16)
$\gamma_{SCS/SF}$	growth rate of SCS/ self-focusing (17/26)
γ_j	ratio of specific heats c_p/c_v for j-type particles (8)
$\Gamma_{e/i/p}$	damping rate of em-wave, ion acoustic wave and plasma wave (12/13/15)
Γ	damping rate (46)
$\Delta\nu$	frequency shift in the back scattered radiation (81)
$\delta\nu$	frequency difference between the two peaks of the spectrum (100)
$\delta\omega$	additional frequency shift due to ion trapping (36)
δ	geometric output coupling parameter (146)
ϵ_0	permittivity of the vacuum (8)
ζ	refractive index of the medium (25)
η	ratio between the electron quiver velocity to the thermal velocity (20)
ϑ	scattering angle between k and k_0 (14)
ϑ_a	first critical angle (147)
ϑ_b	second critical angle (147)
λ_0	laser wavelength (21)
λ_{de}	electron Debye length (8)
μ_0	permeability of the vacuum (8)
ν_{ei}	electron-ion collision frequency (12)
ν_{ii}	ion-ion collision frequency (41)
ν_0	laser frequency (143)
$\varrho(\nu)$	number of photons per unit volume (161)
ϱ_0	laser focal spot radius (29)
τ	relaxation times of the levels (142)

τ_c	life-time of a photon within the laser cavity (162)
ψ_i	ion acoustic wave potential (27)
ψ	electric potential (11)
χ_i	ion susceptibility (17)
	divergence angle (148)
$\omega_{o/\pm}$	frequency of the incident and the back scattered wave (9)
ω	frequency of the ion acoustic wave (8)
ω_{pe}	electron plasma frequency (9)
ω_{pi}	ion plasma frequency (14)
Ω	solid angle (162)

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